

이학박사학위 논문

무거운 강입자를 이용한 입자 생성 연구

Study on the Production of Heavy-Flavor Hadrons

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인하대학교 대학원

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이 논문을 박사학위 논문으로 제출함

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Study on Hadronization Mechanism via Heavy Flavor Hadrons

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Abstract

The ALICE experiment at the LHC at CERN investigated the physics of accelerated heavy-ion collisions at ultra-relativistic energies. These collisions were operated to produce the quark-gluon plasma (QGP) in the laboratory, a deconfined state of the strongly-interacting matter. The QCD theory predicted the formation of QGP under extremely high energy density conditions. Charm and beauty quarks were remarkable probes for investigating the QGP, thanks to their mass order of the GeV. They were created in the hard-scattering processes in the early stages of collision, experiencing the whole evolution of the system. The energy loss of charm and beauty quarks occurred by interacting with the plasma constituents, which can be studied by measuring the heavy-flavor hadrons. In order to separate the effects from the medium and vacuum, a deep comprehension of the hadronization mechanism in pp collisions was required. Recent results on baryon-to-meson production ratio yield in pp collisions at the LHC showed an enhancement with respect to e^+e^- and e^+p collisions. In this thesis, the production cross section of prompt Ξ_c^0 per unit rapidity was described at midrapidity ($|y| < 0.5$) in pp collisions at $\sqrt{s} = 13$ TeV with the ALICE detector. The Ξ_c^0 baryon was measured via the semileptonic decay ($\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e$) channel. The branching-fraction ratio was extracted as $\text{BR}(\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e) / \text{BR}(\Xi_c^0 \rightarrow \Xi^- \pi^+) = 1.38 \pm 0.14 (\text{stat}) \pm 0.22 (\text{syst})$, combining the result of the measurement via the hadronic channel. The total uncertainty of the branching fraction ratio was decreased by a factor of about 3 with respect to the previously published result. The transverse momentum (p_T) dependence of the ratio of Ξ_c^0 and the D^0 , and ratio of the Ξ_c^0 and the $\Sigma_c^{0,+,++}$ and Λ_c^+ were presented. The charm baryon-to-meson ratio increases towards low p_T up to around 0.3. The results provide strong constraints on model calculations and evidence that different charm hadronization processes on taking place in e^+e^- and pp collisions, compared with various hadronization models. ALICE experiment showed outstanding results for the properties of the QGP and the hadronization mechanism. Despite the ambitious scientific program for the upcoming Run 3 and 4, crucial questions still needed to be answered with the present detector. ALICE 3 aimed to investigate the rich phenomenology of QCD matter, parton energy loss, collective flow, hadronization, and electromagnetic radiation. In this thesis, a performance study for multi-charm baryons (Ξ_{cc}^{++}) is discussed to direct the goal of ALICE 3. The significantly higher statistics and innovative detector concept allow the new frontier in heavy-flavor hadron production, providing the new constraint of the charm hadronization mechanism. Quarkonium states, consisting of two pair of heavy quarks, was also used to understand the properties of the QGP. A yield modification of quarkonium states was also observed in small collision systems such as p-Pb collisions, but it is still early to claim any hot medium effect. A model

study in various small collision systems such as pp, p-Pb, p-O, and O-O collisions could provide a quantitatively understand of nuclear effects on the $\Upsilon(nS)$ production. A theoretical calculation considering the gluo-dissociation and inelastic parton scattering and their inverse reaction well described the suppression of $\Upsilon(1S)$ in Pb-Pb collisions. The simulation framework for small systems was studied to investigate the medium effects in small collision systems. In this thesis, the nuclear modification factor of $\Upsilon(nS)$ as a function of charged particle multiplicity ($dN_{\text{ch}}/d\eta$) and transverse momentum are discussed. The elliptic flow of $\Upsilon(nS)$ in small collision systems is presented. All of the results suggest a new reference for upcoming small collisions.



Introduction

Quantum Chromo-Dynamics (QCD) predicts the quark-gluon plasma (QGP), the phase in which quarks and gluons are deconfined within hadrons. The QGP is characterized as the early universe until $\sim 10 \mu\text{s}$ after the Big Bang. The study of the signatures for the QGP can provide a chance to understand the properties of the early universe, which is introduced in Ch. [1](#).

Charm (c) and beauty (b) quarks are mainly produced in the initial hard scatterings, thanks to their high mass ($m_c \simeq 1.3\text{GeV}/c^2, m_b \simeq 4.2 \text{ GeV}/c^2$) [\[49\]](#). Therefore, charm and beauty quark production is a useful tool to probe the QGP, and they survive through the medium expansion, interacting with the free partons in the plasma. The formation of QGP can modify the hadronization mechanisms involving the heavy quarks. However, studying the hadronization process in pp collisions, where the QGP formation is not expected, is mandatory to understand the medium-induced effects thoroughly. In Ch. [2](#), a brief introduction to the open heavy-flavor production in pp collisions, and the quarkonium production in heavy-ion collisions.

The experimental setup will be discussed installed in ALICE (A Large Ion Collider Experiment) in Ch. [3](#), which is one of the four major experiments at the Large Hadron Collider (LHC) of the CERN laboratory in Geneva, Switzerland. For the future experiment, the aims of ALICE 3 are to understand the rich phenomenology of QCD matter using the measurements of multi-charm baryons. The measurements by ALICE in pp collisions provide the constraint of pQCD calculations and for model calculations of charm hadronization in a vacuum. The prompt Ξ_c^0 baryon production decaying into $\Xi_c^0 \rightarrow e\Xi v$ in pp collisions at $\sqrt{s} = 13 \text{ TeV}$ is presented in Ch. [4](#), and the results of the performance study for $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ in pp collisions at $\sqrt{s} = 14 \text{ TeV}$ are introduced in Ch. [5](#).

The response to the medium produced by the collisions affects the physics and yield of the quarkonium state. A model study in different small collision systems, such as pp, p-Pb, p-O, and O-O collisions, will aid in the quantitative understanding of nuclear effects on the production of $\Upsilon(nS)$ and be used to explore the yield modification on $\Upsilon(nS)$ in a small collision system. In Ch. [6](#), the Monte-Carlo framework is introduced In small collision systems to study the medium response on the $\Upsilon(nS)$ states.

Chapter 1

High-energy nuclear physics

The strong interaction between elementary particles, such as quarks and gluons, is explained by Quantum Chromo-Dynamics (QCD), as introduced in Sec. 1.1. The QCD predicts the quark-gluon plasma (QGP), the phase in which quarks and gluons are deconfined within hadrons. The QGP is characterized as the early universe until $\sim 10 \mu s$ after the Big Bang. The QCD also predicts the phase transition of the strongly interacting matter to QGP, as described in Sec. 1.2, where the QCD phase diagram is introduced. The phase transition can happen with the extremely high temperature and high energy density of the system. This condition can be created in ultra-relativistic heavy-ion collisions using the particle accelerator described in Sec. 1.3. The signatures of the QGP will be introduced in Sec. 1.4, with the experimental measurements at the LHC and RHIC colliders.

1.1 Quantum Chromo-Dynamics

One of the fundamental forces in the universe is a strong interaction. It handles nuclear phenomena and explains the interaction between quarks by gluon exchange. In quantum field theory, Quantum Chromo-Dynamics (QCD), which describes the interaction between quarks, is described as a non-abelian gauge theory under a SU(3) color group [50, 51]. The Lagrangian density can be written as:

$$\mathcal{L}_{\text{QCD}} = -\frac{1}{4}G_{\mu\nu}^a G_a^{\mu\nu} + \sum_{f=1}^{n_f} \bar{q}_i^f \left(i\gamma_\mu D_{ij}^\mu - m_f \delta_{ij} \right) q_j^f, \quad (1.1)$$

where q^f is the quark flavor (u, d, c, s, t, b), γ^μ is the Dirac matrices, and δ_{ij} is the Dirac delta function. The term:

$$G_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a - ig_S [A_\mu, A_\nu] = F_{\mu\nu}^a + g_S f_{abc} A_b^\mu A_c^\nu \quad (1.2)$$

indicates the gluon field strength tensor, where a, b, c are color charges in the SU(3) adjoint representation. The gauge field related to the gluon is:

$$A_\mu = \sum_{k=1}^8 A_\mu^k \lambda^k / 2 \quad (1.3)$$

where $A_\mu^a, a = 1, \dots, 8$ are eight gauge fields (conversed charges), and λ^a are the Gell-Mann matrices. The SU(3) algebra is satisfied as defined by the commutation rules, Eq. [1.4](#).

$$\left[\frac{\lambda^a}{2}, \frac{\lambda^b}{2} \right] = i f_{abc} \frac{\lambda^c}{2} \quad (1.4)$$

where f_{abc} the SU(3) structure constants. The first term of Eq. [1.1](#) can be chaged using Eq. [1.2](#):

$$\begin{aligned} \mathcal{L}_G = & \frac{1}{4} F_{\mu\nu}^a F_a^{\mu\nu} \\ & + g_s f_{abc} A_\mu^a A_\nu^b \partial^\mu A_c^\nu - \frac{1}{4} g_s^2 f_{abc} f_{ade} A_b^\mu A_c^\nu A_\mu^d A_\nu^e \end{aligned} \quad (1.5)$$

The Eq. [1.5](#) is QCD Lagrangian, showing the gluonic part. The first line represents the kinetic, and the second line represents the gluon self-interaction which is different from QED. The Eq. [1.5](#) is obtained from the non-abelian group of the SU(3), which shows the additional non-null term of the gluon strength tensor, $-ig_s [A_\mu, A_\nu]$. This characteristic distinguishes QCD and QED. In QCD, gluons have a color charge and can self-interaction, but in QED, the photons have no charge and do not self-interaction.

The covariant derivative is defined as:

$$D_{ij}^\mu = \partial^\mu \delta_{ij} - ig_s \frac{\lambda_{ij}^a}{2} A_a^\mu \quad (1.6)$$

The second term of the QCD lagrangian Eq. [1.1](#) related to quarks can be written using Eq. [1.6](#) as:

$$\begin{aligned} \mathcal{L}_Q = & \sum_{f=1}^{n_f} \bar{q}_i^f \left(i \gamma_\mu D_{ij}^\mu - m_f \delta_{ij} \right) q_j^f \\ = & \sum_{f=1}^{n_f} \bar{q}_i^f \left[i \gamma_\mu \partial^\mu \delta_{ij} - m_f \delta_{ij} \right] q_j^f + g_s \sum_{f=1}^{n_f} \bar{q}_i^f \left[\frac{\lambda_{ij}^a}{2} \gamma_\mu A_a^\mu \right] q_j^f \\ \equiv & \mathcal{L}_0 + \mathcal{L}_I \end{aligned} \quad (1.7)$$

The first part of Eq. [1.7](#), $\mathcal{L}_0$, represents the lagrangian density of the free quark field. This term is invariant under global SU(3) transformations and contains the

quark's kinetic mass information. The second part Eq. [1.7] $\mathcal{L}_I$, is naturally derived by maintaining the lagrangian invariance when local phase SU(3) transformation occurred to $\mathcal{L}_0$ [50].

The adimensional constant g_S indicates the intensity of the strong interaction. This quantity can be described with a strong coupling constant. The strong coupling constant, α_S , depends on the energy scale, and it evolves to the momentum transfer Q according to the equation [52, 53]. The α_S is defined as:

$$\alpha_S(|Q|^2) = \frac{\alpha_S(\mu^2)}{1 + \frac{\alpha_S(\mu^2)}{12\pi} (33 - 2n_f) \ln(|Q|^2/\mu^2)} \quad (1.8)$$

where μ is a renormalization energy scale, and n_f is the number of flavors that contribute to 1-loop corrections, namely those which $m_f < |Q|$ [54].

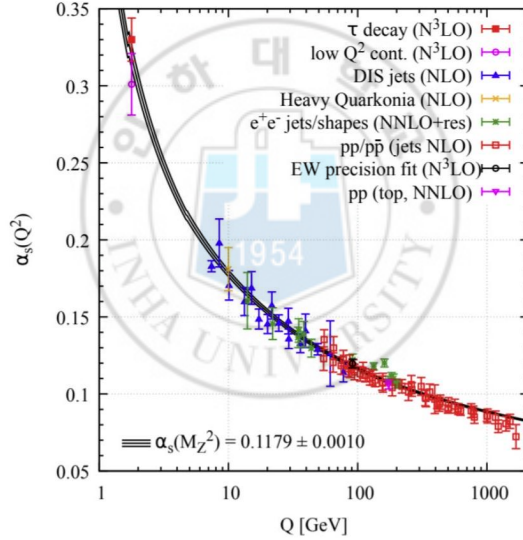


Figure 1.1: The α_S as a function of the momentum transfer Q measured by various experiments. The various perturbative QCD calculations are used to extract α_S . Figure from [1].

The trend of α_S is confirmed by experimental measurement, as shown in Fig. [1.1]. The strong coupling constant rapidly decreases and converges to zero for short distances during the momentum transfer diverges. This trend indicates that the partons can be treated as nonperturbative free particles with infinite energy scales, called asymptotic freedom. Otherwise, the strong coupling constant diverges in large dis-

tances (low energies, $|Q| \lesssim 1 \text{ GeV}/c$), which means that the strong interaction governs higher and higher intensity.

The interaction of a $q\bar{q}$ pair can be treated by the Cornell potential [55]:

$$V(r) = \sigma r - \frac{\alpha}{r} \quad (1.9)$$

where σ is the string tension between the two quarks, and α is a Coulombianlike constant. The equation is valid at $T = 0$.

For a $c\bar{c}$ system, at $T = 0$, the string tension reaches to $\sigma \simeq 0.16 \text{ GeV}^2$ and strong coupling constant reaches to $\alpha \simeq 1/2$ [55]. At high distances, the linear term of the potential increases the system energy until it exceeds the threshold for forming a new $q\bar{q}$ pair, and the $q\bar{q}$ pair can split into two $q\bar{q}$ pairs takes place. The forming of the new $q\bar{q}$ pair in QCD is called confinement. At finite temperatures, the linear term of the potential is tuned to be smaller due to the reduced string tension.

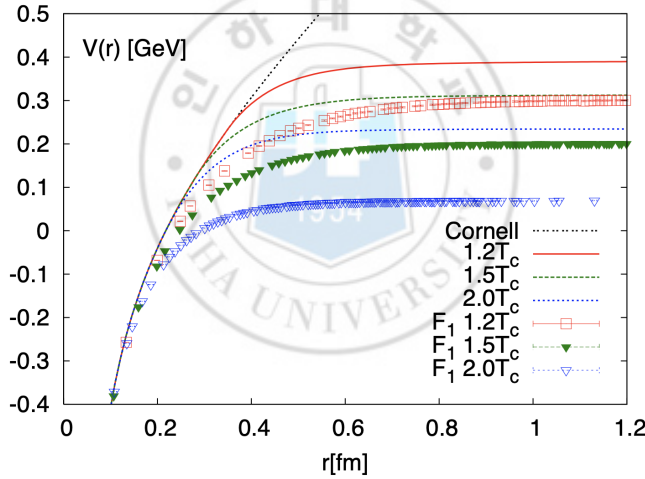


Figure 1.2: Cornell potential (black dash), and confining $q\bar{q}$ potential of heavy quarks calculated with perturbative QCD [2] (color lines), and potential from lattice calculation on singlet free energy [3, 4, 5] (points). Figure from [2].

In Fig. 1.2 Cornell potential is compared with the confining $q\bar{q}$ potential of heavy quarks at various finite temperatures obtained from perturbative QCD [2] and lattice calculation on singlet free energy [3, 4, 5]. For distances larger than $0.4 \div 0.6 \text{ fm}$, the confining potential becomes flat, decreasing with rising temperature. This behavior can explain the suppression of quarkonia in the QGP, as introduced in Sec. 1.4.8.

1.2 Deconfinement and Quark-Gluon Plasma

The α_s constant depends on the energy scale, the theoretical calculations needed to study strong interaction physics are different. The small value of α_s allows perturbative QCD calculations (pQCD) when deriving the scattering amplitude. Different techniques must be adopted if α_s is too large. One of the techniques can be used to evaluate the QCD lagrangian Green's functions under the discrete space-time grid, like lattice QCD (lQCD [56]). The lQCD treats a discrete space-time system corresponding to a complex-time hypercubic lattice occupied by quarks and linked by gluons [56]. The equations from the lQCD are solved in the grid environment. The lQCD predictions can describe the physics phenomena in the limit $a \rightarrow 0$, where a is the spacing between two adjoining grids.

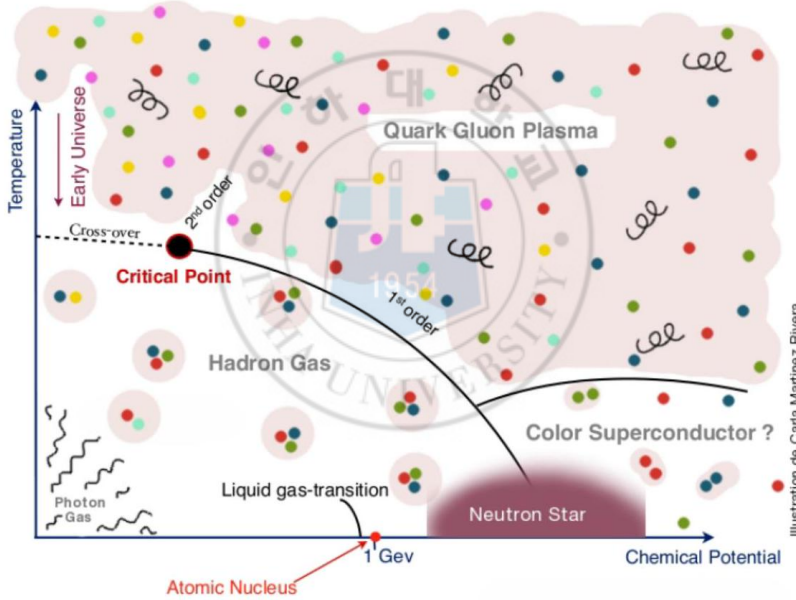


Figure 1.3: The QCD phase diagram. Figure from [6].

Lattice QCD at finite temperature study the point that nearly zero baryon-chemical potential ($\mu_B \approx 0$) on the phase diagram of the strongly-interacting matter, which is described in Fig. [1.3]. The x-axis in Fig. [1.3] means the baryon-chemical potential μ_B , which corresponds to the energy needed to increase the total baryon number by one unity. The μ_B propotional to the net baryon density $n_B - n_{\bar{B}}$. In general, the matter is stable in the region around $T \approx 0$ and $\mu_B \approx 1 \text{ GeV}$. At the high baryon-chemical potential region, which is the high-pressure system above 10^{32} Pa ,

the matter can be turned out as a degenerated gas of neutrons, which is a similar system of the atomic nucleus and could form the core of neutron star [57]. At a higher pressure region, above 10^{35} Pa, the degenerated gas cannot balance for this pressure anymore, and the matter evolves into a low-temperature gas of quarks that are no more captured inside hadrons. Usually, this exotic state of matter is called a color superconductive state [58]. At a high-temperature region, above $T \approx 100\text{MeV}$, the nucleon gas becomes a hadron gas. Moreover, at even higher values, it is expected to transition to the quark-gluon plasma (QGP), a deconfined matter characterized by partonic degrees of freedom. In the low- μ_B region ($\mu_B \approx 0$), the transition to the QGP occurs at a cross-over point which the temperature is $T \approx 154 - 174\text{MeV}$ ($\approx 1.8 - 2.0 \times 10^{12}$ K) and showing the values of the light quark masses $m_{u,d,s}$ [9, 6, 59]. At higher values of the baryon-chemical potential region, the matter experiences a first-order phase transition to QGP, and the co-occurrence with the cross-over at low μ_B indicates the presence of a critical point. Global symmetries of the QCD Lagrangian control the nature of the phase transition. In particular, the cross-over is an outcome of the chiral symmetry restoring to the vanishing quark masses.

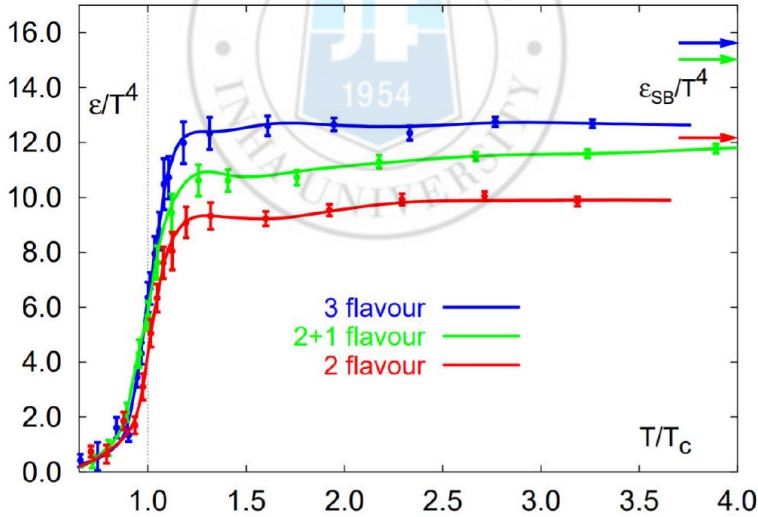


Figure 1.4: Energy density as a function of the temperature of strongly interacting matter at $\mu_B = 0$ from IQCD, with the assumption of either 2 (red) or 3 (blue) light quarks or with 2 light quarks and strange quark (green). The arrows indicate the ϵ/T^4 value predicted for an ideal gas. Figure from [6].

The QGP can be treated as an ideal liquid formed by all elementary particles

with a mass lower than the system temperature by IQCD calculations at finite temperature ($T > \Lambda_{\text{QCD}} \approx 200\text{MeV}$). As shown in Fig. 1.4, the system energy density depends on the fourth power of temperature, $\propto T^4$, after the critical temperature T_c [60]. The ideal gas is considered both bosons and fermions are in chemical equilibrium, and their energy densities follow a Stefan-Boltzmann law:

$$\varepsilon = 3p = g \frac{k_B^4}{h^3 c^3} \frac{\pi^2}{30} T^4, \text{ where } g = \begin{cases} n_{\text{d.o.f.}}, & \text{bosons} \\ \frac{7}{8} n_{\text{d.o.f.}}, & \text{fermions} \end{cases} \quad (1.10)$$

where coefficient g means the number of degrees of freedom from spin, flavors, and charges ($n_{\text{d.o.f.}}$). g has different values for the Bose-Einstein and Fermi-Dirac statistics.

A rapid increase of ε/T^4 is shown in correspondence with the critical temperature T_c that, according to Eq. 1.10, can be described that the number of degrees of freedom is changed while moving from $g = 3$ in the colorless pion gas to the 37 (with two flavors) in the deconfined state [6]. In Fig. 1.4, the trends of 2 light quarks (u, d), 3 light quarks, and 2 light plus 1 strange (s) quarks are shown. The case of 2 light plus 1 strange (s) quarks are closest to the realistic quark mass spectrum and give a critical value of the energy density of $\varepsilon_c \approx 0.7\text{GeV}/\text{fm}^3$ [60].

1.3 Heavy-ion collisions

QGP can be reproduced in heavy-ion collisions at ultra-relativistic energies. The Relativistic Heavy Ion Collider (RHIC) at the Brookhaven National Laboratory (BNL) since 2000 [61] and the Large Hadron Collider (LHC) at CERN since 2010 [62] are the main facilities can generate the ultra-relativistic heavy ion collisions. The maximum center-of-mass energies per nucleon-nucleon pair reach is $\sqrt{s_{\text{NN}}} = 200\text{GeV}$ for gold-gold (Au – Au) collisions at RHIC and $\sqrt{s_{\text{NN}}} = 5.02\text{TeV}$ for lead-lead (Pb – Pb) collisions at the LHC. The net baryon density is very small at those energies since the colliding nuclei are not stopped, allowing us to explore the $\mu_B \approx 0$ area of the QCD phase diagram. This area is called Bjorken, or the transparency regime, different from the stopping area recognized when the accelerated ions are stopped in correspondence with the collision point [63]. The latter area was studied at the Super Proton Synchrotron (SPS [64]) at CERN and the Alternating Gradient Synchrotron (AGS [65]) at BNL, as well as many other facilities during the Beam Energy Scan (BES [66]) at RHIC and the Nuclotron-based Ion Collider Facility (NICA [67]) at the Joint Institute for Nuclear Research in Dubna, Russia.

1.3.1 The Glauber model

The collision geometry is described by the average number of nucleon-nucleon collisions ($\langle N_{\text{coll}} \rangle$) and of participant nucleons ($\langle N_{\text{part}} \rangle$) involved in a particular nucleus-nucleus collision. Protons and neutrons are spread within the nucleus, affecting its dimension since the atomic mass number is related to the nuclear radius by the relation $R \approx 1.2 \text{ fm} \times A^{1/3}$. The Glauber model [7] provides a semi-classical picture of a nucleus-nucleus collision intended as an incoherent superimposition of the compounding nucleons, regarded indistinguishable [7]. In the so-called optical limit, the nucleons move independently in linear trajectories, carrying a momentum to be considered undeflected when the two nuclei intersect. In this way, the collision between two nuclei can be described as a sum of nucleon-nucleon events, where protons and neutrons are spread in the nuclei by Woods-Saxon density ρ . And their interactions are controlled by the inelastic nucleon-nucleon cross section $\sigma_{\text{NN}}^{\text{inel}}$.

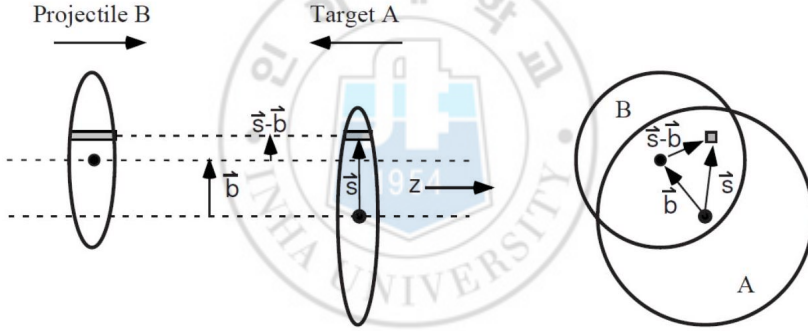


Figure 1.5: Sketch of nucleus-nucleus collision in the optical limit, with transverse (left) and longitudinal (right) views. Figure from [7].

In Fig. 1.5, a sketch of a nucleus-nucleus collision in the optical limit is shown. In the transverse plane, which is orthogonal to the nuclei direction, the system is placed in the center of nucleus A, with respect to the vector $\vec{s}$. The same position for the center of nucleus B is $\vec{b} - \vec{s}$, which is the vector that joins the centers of the two nuclei, called the impact parameter. The $\rho(\vec{s}, z)$, the probability per unit volume to see a nucleon inside the nucleus A found in $\vec{s}$, is determined by the thickness function:

$$\hat{T}_A(\vec{s}_A) = \int \rho_A(\vec{s}_A, z_A) dz_A \quad (1.11)$$

where $\rho_A(\vec{s}_A, z_A) \equiv \rho(\vec{s}, z)/A$. This interpretation is valid for the nucleus B.

The nuclear overlap function is defined as:

$$\hat{T}_{AB}(\vec{b}) = \int \hat{T}_A(\vec{s}) \cdot \hat{T}_B(\vec{s} - \vec{b}) d^2s \quad (1.12)$$

The dimension of $\hat{T}_{AB}(\vec{b})$ is an inverse area, which is related to the probability that a nucleon in A interacts with a nucleon in B. The probability of any nucleon in A colliding with any nucleon in B is described by:

$$p(\vec{b}) = \sigma_{NN}^{\text{inel.}} \cdot \hat{T}_{AB}(\vec{b}) \quad (1.13)$$

where $\sigma_{NN}^{\text{inel.}}$ is the inelastic collisions in nucleon-nucleon collisions. The Eq. 1.13 quantifies the probability of valid single nucleon-nucleon collisions.

If the maximal possible binary collisions between two nuclei with atomic mass numbers A and B is equal to AB, the probability $P_{n,AB}(\vec{b})$ to have n binary collisions out of AB follows the binomial distribution: If the maximal possible binary collisions between two nuclei with atomic mass numbers A and B is equal to AB, the probability $P_{n,AB}(\vec{b})$ to have n binary collisions out of AB follows the binomial distribution:

$$P_{n,AB}(\vec{b}) = \binom{AB}{n} \cdot \left(\sigma_{NN}^{\text{inel.}} \hat{T}_{AB}(\vec{b}) \right)^n \cdot \left(1 - \sigma_{NN}^{\text{inel.}} \hat{T}_{AB}(\vec{b}) \right)^{AB-n} \quad (1.14)$$

where $\binom{AB}{n}$ indicates the possible combinations among nucleons in A and B and $1 - \sigma_{NN}^{\text{inel.}} \hat{T}_{AB}(\vec{b})$ indicates the probability of invalid of single binary collision.

The average number of binary collisions is defined by the first moment of the binomial distribution of Eq. 1.14:

$$\langle N_{\text{coll}} \rangle(\vec{b}) = AB \cdot \sigma_{NN}^{\text{inel.}} \hat{T}_{AB}(\vec{b}) \quad (1.15)$$

The Eq. 1.14 quantifies the total probability of interaction between nuclei A and B.

The interaction probability between A and B can be estimated as the complementary probability of no nucleon-nucleon collision:

$$p_{AB}^{\text{inel.}} = 1 - P_{0,AB}(\vec{b}) = 1 - \left[1 - \sigma_{NN}^{\text{inel.}} \hat{T}_{AB}(\vec{b}) \right]^{AB} \quad (1.16)$$

and the total inelastic nucleus-nucleus cross section can be defined as the probability of the surface integration of the entire transverse plane:

$$\sigma_{AB}^{\text{inel.}} = \int_0^{+\infty} 2\pi b db \cdot \left[1 - \left(1 - \sigma_{NN}^{\text{inel.}} \cdot \hat{T}_{AB}(\vec{b}) \right)^{AB} \right] \quad (1.17)$$

The inelastic nucleus-nucleus cross section in Eq. 1.16 is proportional to the impact parameter.

$$\frac{d\sigma_{AB}^{\text{inel.}}}{db} \propto 2\pi b \quad (1.18)$$

This relation is valid when nuclear is highly overlapped.

Suppose the nuclear overlap region falls when the impact parameter is high and tends to the sum of the colliding nuclei radii. In that case, the linear dependence of inelastic collisions is invalid due to the absence of nucleon-nucleon collisions, given the decreased the participants.

Fixing a nucleon in nucleus A positioned in $\vec{s}$, its probability to collide with a nucleon in B is equivalent with $\sigma_{NN}^{\text{inel.}} \hat{T}_B(\vec{s} - \vec{b})$, where the thickness function is calculated in the position of the nucleon in A, considered with respect to the nucleus B center. The probability of non-interaction with any nucleon in B is $(1 - \sigma_{NN}^{\text{inel.}} \hat{T}_B(\vec{s} - \vec{b}))^B$, same as the product of all failure probabilities. Therefore, the probability for the nucleon in A to be interested in at least one binary collision with a nucleon in B can be defined as $1 - (1 - \sigma_{NN}^{\text{inel.}} \hat{T}_B(\vec{s} - \vec{b}))^B$. Similar to Eq. 1.13, the probability for any nucleon in A to experience at least one binary interaction can be defined as:

$$p_{\text{part}}(\vec{b}) = \int d^2s \cdot \hat{T}_A(\vec{s}) \cdot \left[1 - (1 - \sigma_{NN}^{\text{inel.}} \hat{T}_B(\vec{s} - \vec{b}))^B \right], \quad (1.19)$$

where the integration for all the possible positions of the nucleon in A, whose density per unit area is described by the thickness function $\hat{T}_A(\vec{s})$. Also, the probability of having n participant nucleons out of the A obtainable in the first nucleus is given by a binomial distribution with the number of total cases (A), and the probability of a single success $(1 - (1 - \sigma_{NN}^{\text{inel.}} \hat{T}_B(\vec{s} - \vec{b}))^B)$.

From the first moment of the binomial distribution, the average number of participant nucleons for nucleus A becomes probability in Eq. 1.19 times A. Repeating the same procedure for a nucleon in B, the total number of participant nucleons in A and B can be evaluated by Eq. 1.20.

$$\begin{aligned} \langle N_{\text{part}} \rangle(\vec{b}) = & A \cdot \int d^2s \cdot \hat{T}_A(\vec{s}) \cdot \left[1 - (1 - \sigma_{NN}^{\text{inel.}} \cdot \hat{T}_B(\vec{s} - \vec{b}))^B \right] \\ & + B \cdot \int d^2s \cdot \hat{T}_B(\vec{s} - \vec{b}) \cdot \left[1 - (1 - \sigma_{NN}^{\text{inel.}} \cdot \hat{T}_A(\vec{s}))^A \right] \end{aligned} \quad (1.20)$$

The Monte Carlo (MC) simulation can produce the Glauber model. The nucleons' positions inside the colliding nuclei are specified according to the density

functions ρ , which is the input used to determine the thickness of nuclear (Eq. 1.11). An impact parameter is randomly determined, referring to impact parameter distribution (Eq. 1.17). The interaction between the two nuclei is treated as a superposition of many independent nucleon-nucleon collisions, controlled by the inelastic cross section $\sigma_{NN}^{\text{inel}}$. Two nucleons collide if the distance between them d is less than $\sqrt{\sigma_{NN}^{\text{inel}}}/\pi$.

The MC simulation of the Glauber model is used to determine the average number of participants and binary collisions, for example, at LHC, for central (0 – 10%) Pb – Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV are about $\langle N_{\text{part}} \rangle \simeq 360$ and $\langle N_{\text{coll}} \rangle \simeq 1600$ [68].

1.3.2 Space-time evolution of Heavy-ion collisions

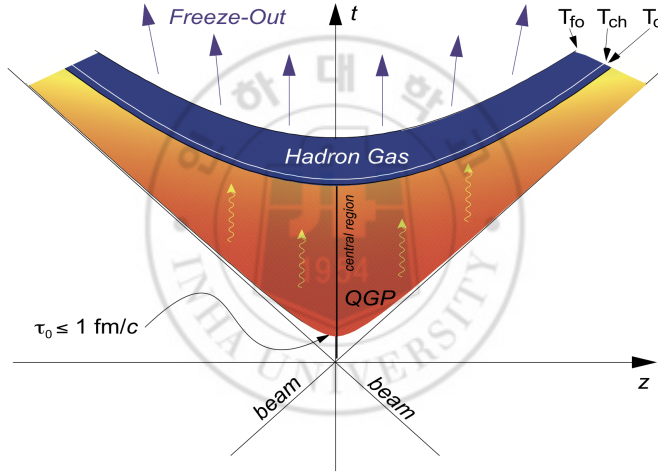


Figure 1.6: Space-time diagram of longitudinal QGP evolution. In ultra-relativistic heavy-ion collisions, the critical energy density is overcome within a time $\tau_0 \sim 1 \text{ fm}/c$, and the QGP phase is reached. The system extends following the hydrodynamics and cools down to the critical temperature T_c , then the transition to the hadron gas phase proceeds. Over time, the particle abundances are confined by the end of inelastic interactions among hadrons at the kinetic freeze-out temperature T_{ch} , followed by a kinetic freeze-out at $T_{\text{fo}} < T_{\text{ch}}$ where the particle energy is definitively established [8]. Figure from [8].

The space-time evolution of the system in a nucleus-nucleus collision is described in Fig. 1.6 and it develops as follows Fig. 1.7.

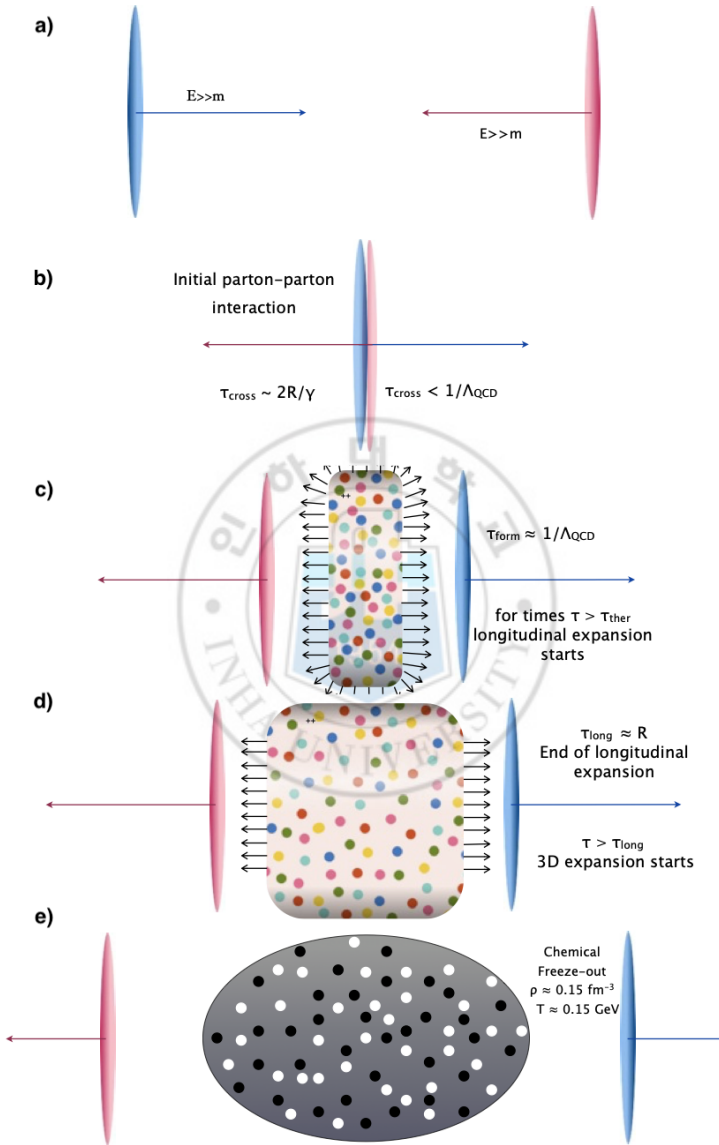


Figure 1.7: Space-time evolution on a heavy ion collision. Figure from [9]

a) Heavy ions are accelerated with ultra-relativistic energies ($E \geq m$) and collide at crossing time τ_{cross} , estimated as:

$$\tau_{\text{cross}} = 2R/\gamma \quad (1.21)$$

where R is the nucleus radius, and γ is the Lorentz factor.

The crossing time for Pb – Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02 \text{ TeV}$ at LHC is known as about $5.7 \times 10^{-3} \text{ fm/c}$, given a Pb radius of about 7.1 fm and a Lorentz factor $\gamma \simeq 2.5 \times 10^3$. The interval τ_{cross} is lower than the strong interaction time scale $\tau_0 = 1/\Lambda_{\text{QCD}} \simeq 1/200 \text{ MeV} \simeq 1 \text{ fm/c}$.

b) The system arrives at a local thermal equilibrium, τ_{ther} , due to multiple scatterings among particles inside the system, characterized by the mean free-path $\lambda \approx 0.1 \text{ fm}$. The expected QGP reproduce temperature at LHC is $\geq 600 \text{ MeV}$ and it decreases as $T \sim \tau^{-1}$ [8]. This stage is considered a preequilibrium state, where the collision's energy density must be sufficient to reproduce the QGP. In the mid-rapidity range, the net baryon density is close to zero. The estimation of system energy density in the Bjorken regime is possible in mid-rapidity range as [63, 8]:

$$\varepsilon = \frac{1}{\pi R^2} \frac{1}{\tau_0} \frac{dE_T}{dy} \Big|_{y=0} \quad (1.22)$$

where the volume produced by initial collisions is considered as a cylinder of length $dz = \tau_0 dy$ and transverse area πR^2 , with $R \approx 1.2 \text{ fm} \times A^{1/3}$, and E_T is the transverse energy.

In RHIC and LHC, at their top center-of-mass energies per nucleon pair with $\tau_0 = 1 \text{ fm/c}$, the energy densities arrived between $\sim 5 \text{ GeV/fm}^3$ and $\sim 12 \text{ GeV/fm}^3$, values significantly higher than the critical value discussed in Sec. 1.2 [69].

c) For $\tau > \tau_{\text{ther}}$, the system evolves by relativistic hydrodynamics laws. First, the longitudinal expansion proceeds because the pressure gradient along the nuclei's direction of motion is more robust than in the transverse plane. Considering the hydrodynamic law [63]:

$$\frac{d\varepsilon}{d\tau} = -\frac{\varepsilon + p}{\tau} \quad (1.23)$$

where p is the system pressure, and ε is the energy density evolved as $\varepsilon \sim 1/\tau^n$, with $1 \leq n \leq 4/3$.

For an ideal ultra-relativistic gas, the component is $n = 4/3$, and the equation of state becomes $\varepsilon = 3p$. This system expansion makes QGP cool down until the time $\tau_{\text{long}} \sim R/c$.

d) For $\tau \geq \tau_{\text{long}}$, the longitudinal system expansion is stopped, and start the tridimensional expansion until the temperature arrives at the critical temperature $T = T_c$. This moment is the phase transition stage from QGP to hadron gas state.

e) After the phase transition, the system reached two subsequent stages:

- Chemical freeze-out: the inelastic interactions among hadrons stop, fixing the species abundances. It happens immediately after the phase transition at $T_{\text{ch}} = 156.5 \pm 1.5 \text{ MeV}$ [70];
- Kinetic freeze-out: for $T_{\text{fo}} < T_{\text{ch}}$ with also elastic interactions among hadrons stops, energy and momentum of the final particles are defined.

1.4 Signature of Quark-Gluon Plasma

This section reports an overview of the experimental results and the experimental observables that mainly contribute to the fundamental understanding of the QGP properties.

1.4.1 Chemical freeze-out temperature

At the end of the evolution of the QGP, the hadronization occurred from partons into different hadronic species. Soft hadrons production is impossible in a microscopic view due to the channels that may contribute to the non-perturbative process and the rescattering of soft partons in-medium that obliterates their initial production. Therefore, a statistical approach is introduced, that the hadron is produced macroscopically in the system [71]. At chemical freeze-out temperature, the system is assumed to reach the thermal and chemical equilibrium, corresponding to a hadron-resonance gas forming an ensemble for conservation of energy and quantum numbers [70]. Under these assumptions, the system thermally produced the abundance N_i of a given particle species i with volume V . Furthermore, temperature T_{ch} is obtained with the grand-canonical ensemble formalism:

$$N_i(T_{\text{ch}}, V, \mu_i) = \frac{g_i V}{(2\pi)^3} \int \frac{d^3 p}{\exp\left(\frac{E_i - \mu_i}{T_{\text{ch}}}\right) \pm 1} \quad (1.24)$$

where g_i is the spin degeneracy, E_i is the particle's energy, and the integral is completed in the entire momentum space. The " $\pm$ " symbol considers the Bose-Einstein and Fermi-Dirac statistics, μ_i is the chemical potential for each particle governed by

quantum numbers, $\mu_i = \mu_B B_i + \mu_S S_i$, where B_i is the baryonic number and S_i the strangeness number.

The Eq. 1.24 may not be formally proper in a non-perfect thermal and chemical equilibrium. For example, the behavior of a strange quark may differ from the expectation by the grand canonical description. In this case, the thermal abundance in Eq. 1.24 is scaled by the strangeness suppression factor, $\gamma_S^{\delta_i}$. γ_S is expected to be lower than unity in non-equilibrium, such as peripheral AA collisions and small collision systems. Similar statements can be interpreted for the charm quark (γ_c), mainly produced in the initial hard scatterings. Finally, another parameter γ_q is introduced in the SHARE model, introducing the hadronization without additional re-scatterings in non-equilibrium QGP.

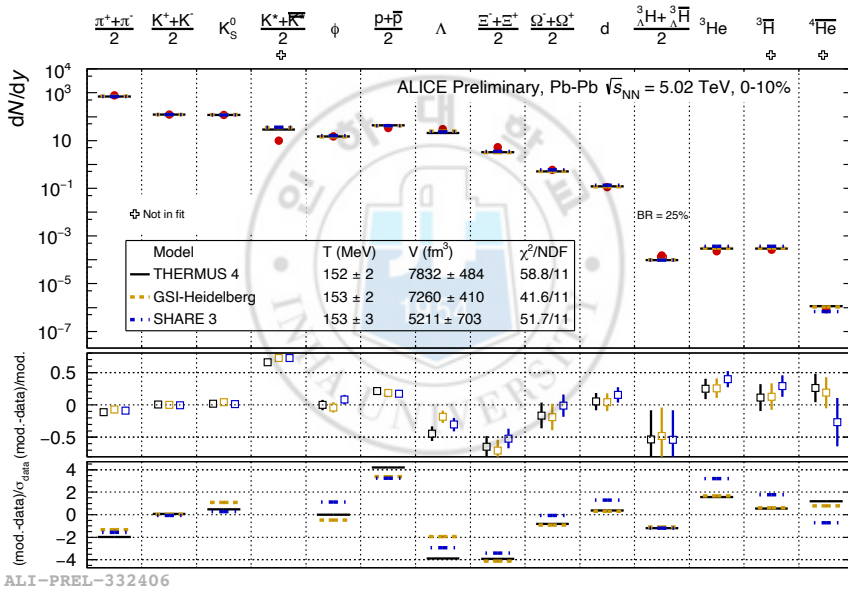


Figure 1.8: Simultaneous fit of particle yield measured by ALICE experiment in 0 – 10% Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. Figure from ALICE Figure Repository.

The particle yields per unit of rapidity measured by the ALICE in 0 – 10% Pb – Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV are shown in Fig. 1.8. The experimental data are compared with the model calculations of SHARE [72], THERMUS [73], and GSI [74], which assumes the different equilibrium condition of the system, and/or conservation laws at the chemical freeze-out. And they fit the different particle yields leaving T_{ch} , V , and μ as free parameters. The model calculations explain the particle abundances in a magnitude of about 9 orders with a χ^2/NDF be-

tween 3.8 and 5.3. The measured yield of excited K mesons and proton-antiproton are excluded from the fits due to the sizable annihilation cross section of baryon-antibaryon in the hadronic phase [75]. All the models show the chemical freeze-out temperature is about $T_{\text{ch}} = 153 \text{ MeV}$. This result is close to the expected critical temperature of the phase transition from the hadronic to the deconfined state, indicating that the chemical freeze-out should happen very early stage [72, 73, 74].

1.4.2 Kinetic freeze-out temperature

After the chemical freeze-out, the particle abundances are set, but the particle's energy can not be fixed due to residual elastic interactions. The particle p_T spectra are fixed after the kinetic freeze-out. Low p_T region can give a fundamental understanding of the medium properties at the stage of evolution. In thermal equilibrium, the low- p_T spectrum of particles follows a Boltzmann-Gibbs law:

$$\frac{1}{p_T} \frac{d^2 N}{dp_T dy} = \frac{1}{m_T} \frac{d^2 N}{dm_T dy} \propto e^{-m_T/T_{\text{kin}}} \quad (1.25)$$

where m_T is the transverse mass ($m_T := \sqrt{m^2 + p_T^2}$), and T_{kin} is kinetic freeze-out temperature.

According to Eq. 1.25, all particle species follow the exact exponential trend, and this trend is usually called m_T -scaling [10]. The different particle species (e.g. π^- , $\bar{K}^0$, Λ , ρ^0 and ω) produced in hadronic collisions at 8 and 16 GeV [76] are well described by Eq. 1.25. The measurement in pp collisions at 7 TeV with ALICE [77] challenges the universal validity. The hadron p_T spectra deviate from this trend at the LHC energies. This deviation is also observed in Pb – Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76 \text{ TeV}$ [78] and 5.02 TeV [10]. The reason for disagreement is due to the pressure gradients caused by the QGP, which produce a collective flow in the transverse plane, called radial flow, that superimposes the random thermal motion of particles. Therefore, the slope parameter in Eq. 1.25 needs to be changed as Eq. 1.26.

$$\bar{T} = T_{\text{kin}} + \frac{m}{2} \langle \beta_T \rangle^2 \quad (1.26)$$

where the second term indicates the kinetic energy related to the radial expansion velocity $\langle \beta_T \rangle$, which forces up to higher values the average p_T of the particles, proportionally to their mass. β_T is not constant, but transformations with a power law of the distance r for the center of the expanding system:

$$\beta_T(r) = \beta_s \left(\frac{r}{R} \right)^n \quad (1.27)$$

where β_s is the surface velocity of the expanding system, and R is the maximum radius of the fireball.

The particle production as a function of momentum is introduced by the Boltzmann-Gibbs Blast-Wave function [79], which contains T_{kin} , $\langle\beta_T\rangle$ and n as free parameters. This method allows for distinguishing the random thermal motion from the possible collective motions from the measured spectra of the different hadron species. In the Blast-Wave model, each microscopic volume element of the fireball is considered to reach the thermal equilibrium at a specific temperature T and experiences a radial boost with velocity β_T . The spectrum is finally affected by a superposition of thermal origins with different velocities of radial expansion.

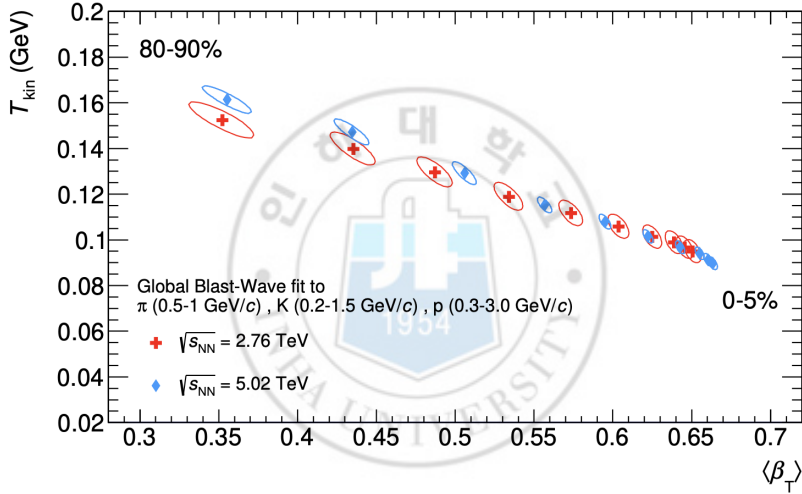


Figure 1.9: Average expansion velocity ($\langle\beta_T\rangle$) and kinetic freeze-out temperature (T_{kin}) from the Boltzmann-Gibbs Blast-Wave fit to $\pi^\pm$, $K^\pm$ and $p(\bar{p})$ spectra measured in Pb – Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ and 2.76 TeV. The correlated uncertainties of the global fit are shown as ellipses. The elliptic outlines correspond to 1σ uncertainties, with statistical and systematic uncertainties added in quadrature. Figure from [10].

The correlation between the average expansion velocity β_T and the kinetic freeze-out temperature T_{kin} measured by the ALICE experiment in Pb – Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [78] and 5.02 TeV [10] are shown in Fig. 1.9. The two parameters are estimated with a simultaneous fit of pion, kaon, and proton p_T spectra at low p_T measured at different centralities. For Pb – Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV, the

average transverse expansion arrives at a maximum value of $\langle\beta_T\rangle = 0.663 \pm 0.003$ in 0–5% centrality class, while the kinetic temperature decreases to central collisions, varying between $T_{\text{kin}} = (0.161 \pm 0.006)$ GeV and $T_{\text{kin}} = (0.090 \pm 0.003)$ GeV [10]. These results suggest a more rapid expansion with increasing centrality, given the more extensive system size. They also show that the kinetic freeze-out is delayed in central collisions because the system cools down at lower temperatures. No significant collision energy dependence is observed. The values of T and $\langle\beta_T\rangle$ calculated from the simultaneous fit show a physical meaning only if model calculations consider the dynamical expansion of the system.

1.4.3 Electromagnetic emissivity

Direct photons are useful for studying the time evolution of the QGP produced in heavy-ion collisions because they are produced at all collision stages and traverse the medium individually since they do not interact strongly with the medium [80]. As a result, their mean-free path is much larger than nuclear scales. The prompt direct photons produced in the hard scatterings among partons in the colliding nuclei provide information on the parton densities in nuclei [81]. The prompt direct photons dominate the high- p_T distribution of the spectrum ($p_T > 5$ GeV/ c). On the other hand, the thermal photons dominate the low- p_T region [11].

The direct-photon p_T spectra in central (0-20%), semicentral (20-40%) and peripheral (40-80%) Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV measured by ALICE are shown in the left panel of Fig. 1.10. The pQCD calculations well describe the direct photon distribution for $p_T > 5$ GeV/ c at the same energy scaled by the expected $\langle N_{\text{coll}} \rangle$ for each centrality class. The power-law dependency of p_T , predicted from the hard process, is broken at low p_T due to an overabundance of thermal photons having an exponential law dependency in semicentral and more clearly in central collisions.

The parameter of inverse slope for the exponential shape indicates an effective temperature of the medium (T_{eff}) averaged over an entire evolution, according to the theoretical models. A fit of the ALICE measurement in $0.9 < p_T < 2.1$ GeV/ c region with the exponential function $A \cdot \exp(-p_T/T_{\text{eff}})$ provides an effective temperature of the QGP produced in central Pb–Pb collisions, which at $\sqrt{s_{\text{NN}}} = 2.76$ TeV corresponds to $T_{\text{eff}} = (304 \pm 11 \text{ (stat.)} \pm 40 \text{ (syst.)})$ MeV [11]. The effective temperature is 1.8 times higher than the critical temperature T_c (see Sec. 1.2). The QGP produced in central Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV at LHC is about 1.3 times higher than one in Au–Au collisions at $\sqrt{s_{\text{NN}}} = 200$ GeV in the same centrality class at RHIC [11].

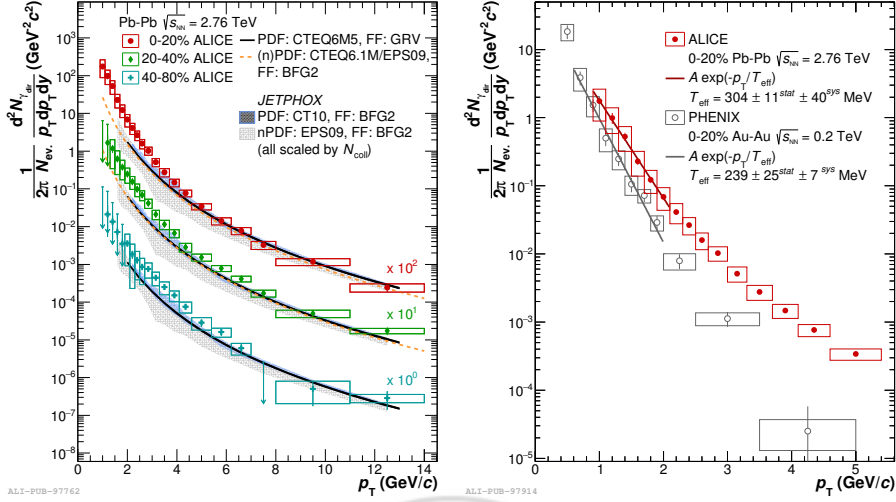


Figure 1.10: Left: direct-photon spectrum measured by the ALICE in Pb – Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV for the 0 – 20% (scaled by a factor 100), the 20 – 40% (scaled by a factor 10) and 40 – 80% centrality classes. NLO pQCD predictions are scaled by the number of binary nucleon collisions for each centrality class for the direct photon yield in pp collisions at the same energy. Right: direct photon spectra in Pb – Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV and Au – Au collisions at $\sqrt{s_{NN}} = 200$ GeV measured by the ALICE at the LHC and the PHENIX at RHIC. For $p_T < 2.1$ GeV/c the exponential function $A \cdot \exp(-p_T/T_{\text{eff}})$ is used to derive the average temperature of the medium. Figures from [11].

1.4.4 Charged particle multiplicity and energy density

The number of charged particles produced in a collision can be described with the charged particle multiplicity. The charged particle multiplicity corresponds to the energy density affected in the pre-equilibrium phase described in Sec. 1.3.2, which decides whether the matter experiences a transition to the QGP. The energy density at midrapidity reported in Eq. 1.22 can be approximated as:

$$\left. \frac{dE_T}{dy} \right|_{y=0} \approx \langle E_T \rangle \cdot \left. \frac{dN}{dy} \right|_{y=0} \quad (1.28)$$

where the last component corresponds to the rapidity distribution of charged particles at midrapidity. The pseudorapidity is interpreted as:

$$\eta := -\ln \left(\tan \frac{\theta}{2} \right) \quad (1.29)$$

where θ is the polar angle, defined in the $y-z$ plane (see Sec. 3.2).

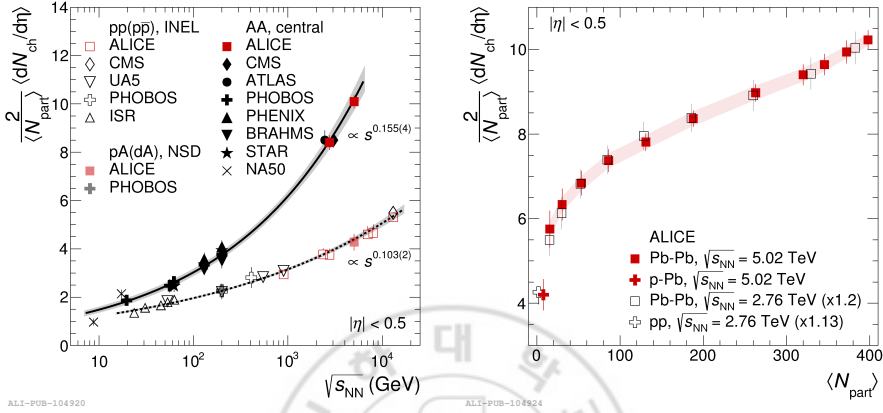


Figure 1.11: Left: Charged particle pseudorapidity densities per participant nucleon pair $2/\langle N_{\text{part}} \rangle \langle dN/d\eta \rangle$ measured at mid rapidity in different collision systems as a function of center-of-mass energy per nucleon pair $\sqrt{s_{\text{NN}}}$. Right: $2/\langle N_{\text{part}} \rangle \langle dN/d\eta \rangle$ as a function of number of participants N_{part} . Figures from [12].

The pseudorapidity density of the charged particle at mid rapidity in pp, p – Pb and Pb – Pb collisions measured by ALICE is shown in the left panel of Fig. 1.11. The pseudorapidity density of the charged particle is normalized by the number of participant pairs $\langle N_{\text{part}} \rangle / 2$ to compare among different collision systems [12]. The charged particle pseudorapidity density per participant nucleon pair $2/\langle N_{\text{part}} \rangle \langle dN/d\eta \rangle$ measured by ALICE is compared with the results from other experiments at different $\sqrt{s_{\text{NN}}}$. The particle multiplicity shows a monotonic growth as a function of $\sqrt{s_{\text{NN}}}$. The results from central AA collisions follow a power law function as^b with parameter b equal to 0.155 ± 0.004 , while the results from pp collisions following a power law increase regulated by $b = 0.103 \pm 0.002$ [12]. The results in pA and dA collisions lie on the pp curve, suggesting that the trend observed in AA collision cannot be attributed only to a more significant number of multiple interactions among the participant nucleons, also present in pA and dA collisions [12].

In the right panel of Fig. 1.11, the $2/\langle N_{\text{part}} \rangle \langle dN/d\eta \rangle$ shows a significant enhancement as a function of the number of participant nucleons. The charged particle multiplicity in Pb – Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02 \text{ TeV}$ increases almost linearly

with $\langle N_{\text{part}} \rangle$ and significantly differs from measured in pp and p–Pb collisions at the same energy. This suggests that the charged particle multiplicity does not depend on the collision energy. Instead, it is described by models that explain the lower multiplicity in smaller systems to the saturation of parton densities, which limits the production of partons and, finally, of particles [12].

1.4.5 Strangeness enhancement

The mass of the strange quark is $\sim 0.5 \text{ GeV}/c^2$, suggesting a production threshold for a $s\bar{s}$ pair larger than 1 GeV, below the critical temperature T_c . Thanks to the chiral symmetry restoration in QGP (see Sec. 1.2), above the critical temperature, the effective mass of the strange quark reduces by a factor up to 10, resulting increase of the $s\bar{s}$ pair production [82]. For this reason, the strangeness production is predicted to increase in heavy-ion collisions compared with smaller systems. The strangeness enhancement was proposed as a signature of the QGP formation in heavy-ion collisions [78]. Strange quarks are expected to be abundantly produced in the QGP via the gluon fusion, $gg \rightarrow s\bar{s}$, thanks to the huge gluon density in the system [83].

The strangeness enhancement in heavy-ion collisions can be demonstrated in a reverse effect as strangeness suppression in small collision systems [79]. The conservation laws in QGP need to be respected inside the system's volume, according to the grand-canonical ensemble. The conservation is not required locally but only on average in the entire volume, controlled by the related chemical potential. This means that the strange hadron production at a specific position in the fireball does not need the production of anti-strange hadron nearby, but this can occur everywhere inside the system volume.

On the other hand, in small collision systems, such as pp, the grand-canonical ensemble cannot be used because the reduced dimensions of the phase space need a canonical ensemble. In small collision systems, the conservation of all the quantum numbers must take place locally. This means that the strange hadron production denotes the anti-strange hadron production in the same space-time position. For these reasons, the strangeness enhancement in the QGP is usually diagnosed as the reduction of the so-called canonical suppression in small collision systems. This is expressed as a γ_s factor lower than unity in pp collisions (see Eq. 1.24), indicating the deviation from the grand-canonical ensemble valid in heavy-ion collisions ($\gamma_s = 1$).

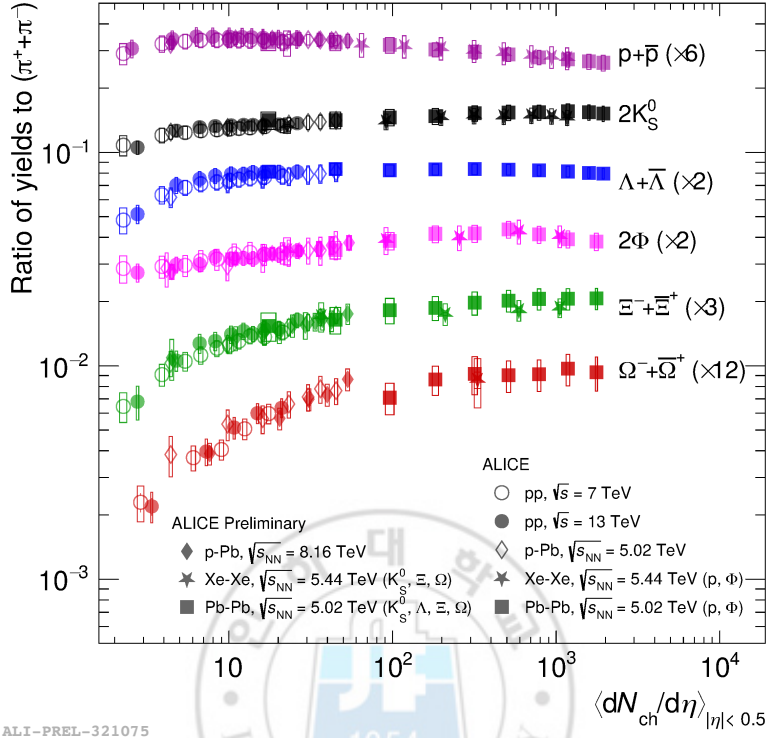


Figure 1.12: Relative yields to $(\pi^+ + \pi^-)$ of $(p + \bar{p})$, $2 K_S^0$, $(\Lambda + \bar{\Lambda})$, 2Φ , $(\Xi^- + \bar{\Xi}^+)$ and $(\Omega^- + \bar{\Omega}^+)$ as a function of charged particle multiplicity at midrapidity measured by the ALICE in pp collisions at $\sqrt{s} = 7$ and 13 TeV, in p – Pb collisions at $\sqrt{s_{NN}} = 5.02$ and 8.16 TeV, in Pb – Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV and in Xe–Xe collisions at $\sqrt{s_{NN}} = 5.44$ TeV. Figure from ALICE Figure Repository.

Strangeness enhancement was observed in heavy-ion collisions at RHIC [83] and LHC [84]. The production yields normalized to those of $(\pi^+ + \pi^-)$ of $(p + \bar{p})$, $2 K_S^0$, $(\Lambda + \bar{\Lambda})$, 2Φ , $(\Xi^- + \bar{\Xi}^+)$ and $(\Omega^- + \bar{\Omega}^+)$ as a function of charged particle multiplicity at mid rapidity $\langle dN/d\eta \rangle_{|y| < 0.5}$ measured by the ALICE in pp collisions at $\sqrt{s} = 7$ and 13 TeV, in p – Pb collisions at $\sqrt{s_{NN}} = 5.02$ and 8.16 TeV, in Pb – Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV and in Xe–Xe collisions at $\sqrt{s_{NN}} = 5.44$ TeV are shown in Fig. 1.12. The measurements show increased production of strange hadrons in Pb – Pb collisions with respect to the low-multiplicity pp events, showing a more significant enhancement for multi-strange-quark particles. The Ω^- (sss) baryons are highest for the relative increase between the two collision systems, whose production increases by a factor ~ 5 in Pb – Pb compared

with low-multiplicity pp collisions. The measurements do not show a significant dependence on the center-of-mass energy per nucleon-nucleon pair or the collision system, showing compatible production yields in collisions producing similar final state multiplicity $\langle dN/d\eta \rangle|_{|y|<0.5}$ [84]. In particular, recent results from the ALICE [85] show that the strange hadron production in high-multiplicity pp collisions is comparable with those measured in Pb–Pb collisions.

1.4.6 Azimuthal anisotropy

The collision geometry of a nucleus-nucleus collision depends on its impact parameter, as described in Sec. 1.3.1. The event is defined as central when the impact parameter is close to zero. The event is defined as noncentral when the impact parameter is larger, and only a fraction of the nucleons experience the interaction while the others move forward unaffected.

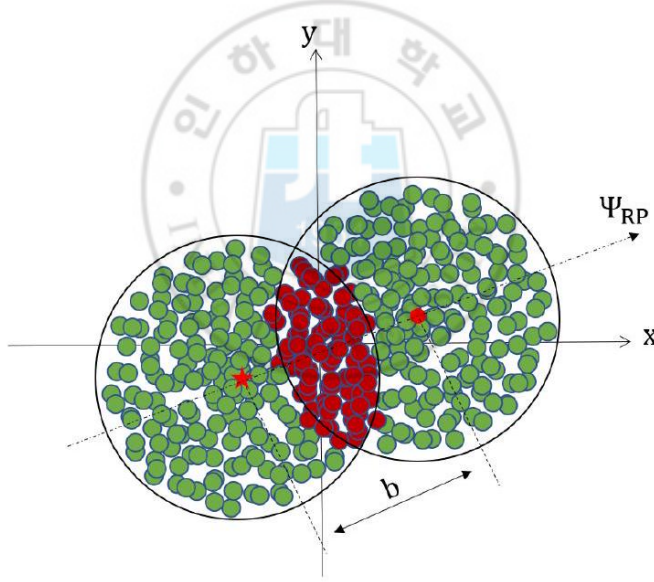


Figure 1.13: The collision between two lead ions. The direction of motion $\hat{z}$ is perpendicular to the transverse plane xy . The quantity Ψ_{RP} identifies the direction of the reaction plane. The impact parameter b is defined as the distance in the transverse plane between the ion centers.

As shown in Fig. 1.13, in non-central events, the overlapping area identified by the participant nucleons assumes an almond-like shape. The projection of overlapping areas in the transverse plane is like an ellipse shape. This geometrical

anisotropy of the initial state is transformed into an anisotropy in momentum space due to the anisotropic pressure gradients in the transverse plane [86]. The anisotropy in momentum space causes a parton's azimuthal anisotropy in the collective flow before their hadronization [86]. The measurement of the anisotropic flow enables the study of the equation of state of the system and the transport coefficients related to the hydrodynamic expansion, such as the shear viscosity to entropy density η/s and the bulk viscosity to entropy density ζ/s [87, 88].

The azimuthal anisotropy of the system can be interpreted through the measurement of the hadron production as a function of the azimuthal angle ϕ , which can be decomposed with a Fourier expansion [89] as:

$$E \frac{d^3N}{dp^3} = \frac{1}{2\pi} \frac{d^2N}{p_T dp_T dy} \left(1 + 2 \sum_{n=1}^{\infty} v_n(p_T) \cos[n(\phi - \Psi_n)] \right) \quad (1.30)$$

where Ψ_n in Eq. [1.30] is the symmetry plane of the harmonic of n -th order.

The quantity v_n indicates the n -th order Fourier coefficient and is calculated as:

$$v_n(p_T) = \langle \cos[n(\phi - \Psi_n)] \rangle \quad (1.31)$$

where the brackets are the average among all the particles with a given p_T in all the events with a given centrality.

The v_2 is called elliptic flow and constitutes the significant contribution to the asymmetry in a non-central collision by the almond-like shape. The measurement of a large v_2 is evidence of the QGP, indicating the presence of a thermalized system that follows the laws of relativistic hydrodynamics already in the partonic phase [90]. The higher order coefficients v_3 and v_4 , called respectively triangular and quadrangular flow coefficients, are generated mainly by fluctuations of the distributions of participants in the nuclei in the initial state and by the non-linear hydrodynamic response of the medium [91, 86, 90, 92, 93, 94].

Studying the anisotropic flow for identified particles suggests a more profound knowledge of the QGP evolution's physics. In particular, the p_T dependence allows studying different kinematic regions of interest, where different physics phenomena affect the flow. The p_T -differential elliptic flow coefficient of $\pi^\pm, K^\pm, \phi, p + \bar{p}, \Lambda + \bar{\Lambda}$ and K_S^0 measured by the ALICE experiment in Pb – Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV at different centrality classes is shown in Fig. [1.14] [13]. The v_2 increases from most central collisions (0 – 1%) to semi-central ones (30 – 40%) because of the increased the overlapping region between the two colliding nuclei. For more peripheral collisions, the v_2 decreases again due to the lower pressure gradients in the created system. In the low- p_T region, ($p_T \lesssim 3$ GeV/ c), the elliptic flow shows a clear mass ordering, namely the v_2 coefficient decreases with the increasing hadron

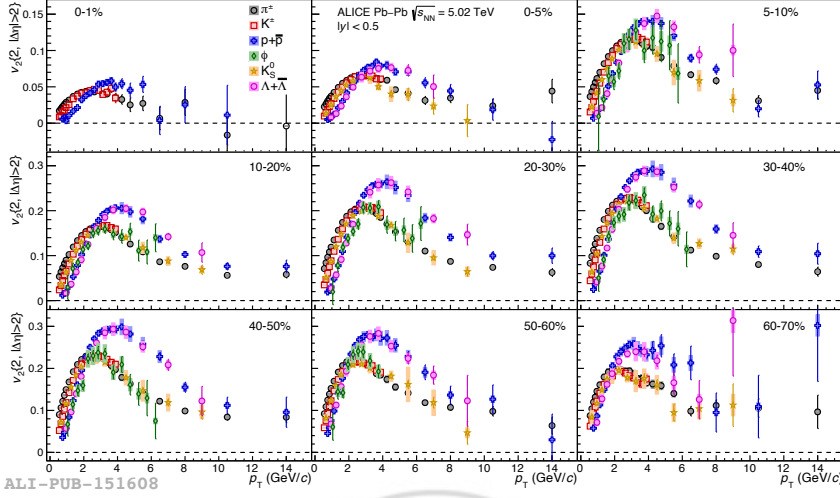


Figure 1.14: p_T -differential v_2 of $\pi^\pm, K^\pm, p + \bar{p}, \Lambda + \bar{\Lambda}, K_S^0$ and ϕ meson in Pb – Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV at different centrality classes. Figure from [13].

mass. This behavior is expected due to the interplay with the radial flow (see Sec. 1.4.2) pushing more massive particles to higher $\langle p_T \rangle$ [95]. At intermediate p_T region, ($3 \lesssim p_T \lesssim 8 \text{ GeV}/c$), v_2 tend to be higher for baryons with respect to mesons. This property supports the quark hadronization via the coalescence, which implies that the v_2 of each single constituent quark contributes to the elliptic flow of the final state hadron. This implies a larger v_2 for baryons with respect to mesons, showing a v_2 dependence on the number of constituent quarks rather than the hadron mass [95]. The measurement of anisotropic flow at high p_T , when $p_T \gg m$, is sensitive to the in-medium path-length dependent energy loss effects. For transverse momenta $p_T \gtrsim 8 \text{ GeV}/c$ the anisotropic flow is procured when hard partons transport through the QGP losing energy via multiple scatterings and gluon radiation, leading to a non-zero elliptic flow up to very high p_T [94, 96, 97, 98].

The experimental results support the description of the QGP as a perfect fluid since they are reasonably described by model calculations [99, 95], which assume a η/s close to the lower limit of $1/4\pi$. This value is imposed by first principle calculations and confirmed by anti-de Sitter/conformal field theory (AdS/CFT [87]), even if recent theoretical developments seem to challenge this limit [88].

1.4.7 Jet quenching

The production of high-energy partons in hadronic collisions at leading order produces a pair of back-to-back high-energy partons. A jet is a spray of hadrons emanating from the parton shower generated by the high-energy parton fragmentation. The back-to-back topology and the kinematic correlations are broken in nucleus-nucleus collisions due to the QGP. As shown in Fig. 1.15, the two partons produced in the hard-scattering process lose energy as gluon radiation traverses the medium, where energy loss depends on the path-length. This phenomenon is called jet quenching and was observed at RHIC [15], comparing the di-hadron azimuthal correlation in nucleus-nucleus collisions with the results from small collision systems.

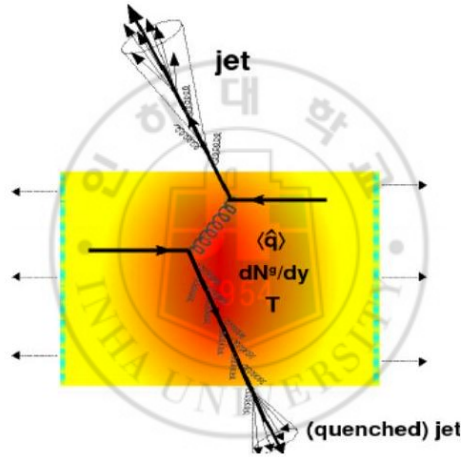


Figure 1.15: Sketch of jet quenching in QGP. Figure from [14].

Quarks produced close to the fireball surface preserve most of their original energy and are later observable as high- p_T jets. This is usually used as a possible "trigger" particle, and the region with $\Delta\phi = 0$ is called near-side. The opposite parton, expected in the away-side region, at $\Delta\phi = \pi$ with respect to the trigger particle, must cover a long path before leaving the QGP, losing more energy.

The measurements from the STAR in pp, deuteron-gold (d – Au) and Au – Au collisions at $\sqrt{s_{NN}} = 200$ GeV are shown in Fig. 1.16. While the near-side peak has a comparable height and shape among different collision systems, the away-side peak in central (0-20%) Au-Au collisions vanishes, signaling that the back-scattered parton loses scattered energy at large angles. This behavior is not observed in d-Au collisions, where no medium formation is expected.

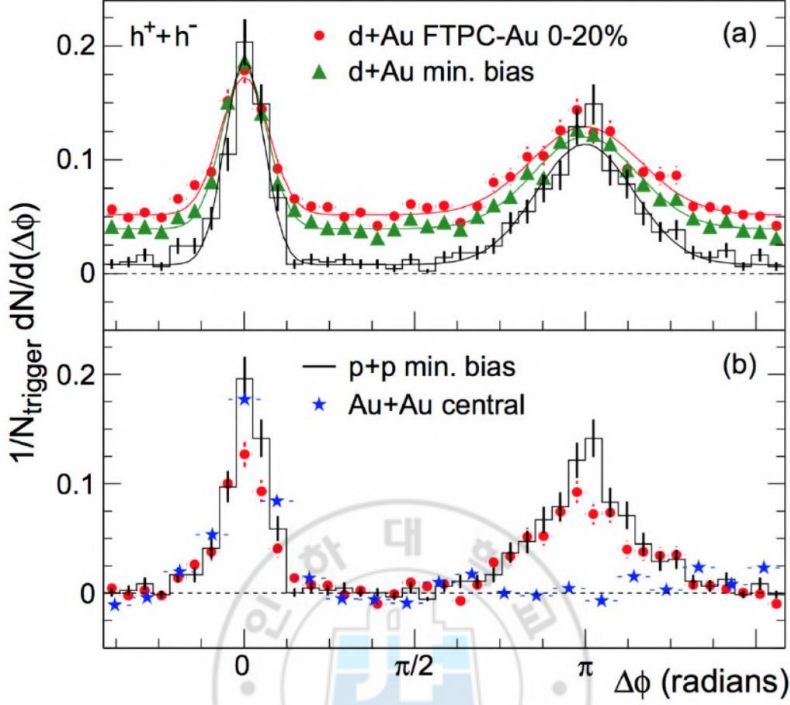


Figure 1.16: (a) azimuthal distributions for minimum bias and central d-Au collisions, and pp collisions at $\sqrt{s_{\text{NN}}} = 200 \text{ GeV}$ measured by the STAR at RHIC. (b) comparison of pedestal subtracted two-particle azimuthal distributions for central d-Au collisions to those seen in pp and central Au – Au collisions at $\sqrt{s_{\text{NN}}} = 200 \text{ GeV}$ measured by the STAR at RHIC. Figure from [15].

The hadrons possibly produced from the away-side are influenced by the loosing of the quark energy, inducing an all-around softening of the measured p_T -spectrum of final state particles. This effect is studied with the measurement of the nuclear modification factor (R_{AA}), which is defined as follows:

$$R_{\text{AA}}(p_T, y) = \frac{1}{\langle T_{\text{AA}} \rangle} \frac{d^2 N_{\text{AA}} / dp_T dy}{d^2 \sigma_{\text{pp}} / dp_T dy} \quad (1.32)$$

where $\langle T_{\text{AA}} \rangle = \langle N_{\text{coll}} \rangle / \sigma_{\text{pp}}^{\text{inel.}}$ (see Eq. 1.15). The Eq. 1.32 is described with the p_T - and rapidity-differential yield measured in nucleus-nucleus collisions divided by in pp collisions scaled by the average number of expected nucleon-nucleon interactions, from the Glauber model (see Sec. 1.3.1). The nuclear modification factor is expected to be unity in case of no QGP scenario.

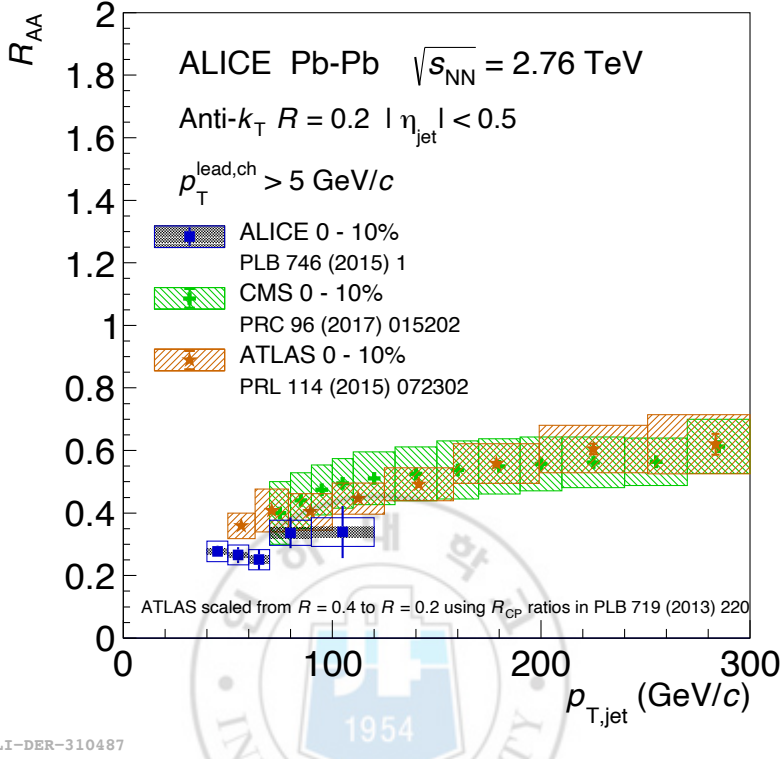


Figure 1.17: Nuclear modification factor of charged jets in central (0–10%)Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV measured by the ALICE experiment [16] compared with one of inclusive jets measured by the CMS [17] and ATLAS [18] experiments at the same collision energy and centrality. Figure from ALICE Figure Repository.

The charged particle R_{AA} of all centrality classes increases for $p_T \gtrsim 10$ GeV/c is observed toward unity, a fairly constant p_T dependence is observed for charged jets within the same p_T interval. These trend is shown in Fig. 1.17, where the nuclear modification factor of charged jets in central (0–10%)Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV measured by the ALICE [16] compared with inclusive jets measured by CMS [17] and ATLAS [18]. The trend of increase R_{AA} of charged particles is not accompanied by a similar trend of the charged jet R_{AA} , which reaches a plateau at around 0.5 for $p_{T,jets} \gtrsim 200$ GeV/c. This comparison suggests that part of the energy is spread out of the jet cone in Pb–Pb collisions, causing a R_{AA} for charged jets lower than unity.

1.4.8 Quarkonia suppression

The study of $c\bar{c}$ and $b\bar{b}$ bound states, called quarkonia, is useful to inspect the QGP properties. The suppression of the charmonia production in heavy-ion collisions is sensitive for deconfinement and sensitive to the medium temperature [2]. The quarkonium binding energy in the vacuum can be derived from the Cornell potential reported in Eq. 1.9, which allows to define the bound state radius that for the J/ψ ranges in $0.2 \leq r_{J/\psi} \leq 0.5\text{fm}$ at $T = 0$, depending on the charm quark mass and of the string tension σ and α at zero temperature. At finite T , the string tension coefficient decreases until $\sigma(T_c) = 0$, the linear confining term of the potential vanishes for $T \geq T_c$, when it becomes:

$$V(r)^{\text{QGP}} = -\frac{\alpha(T)}{r} \cdot e^{-r/r_D(T)} \quad (1.33)$$

The Cornell potential is reduced to the color screen potential, and the color charge obtains a screening effect similar to that experienced by the charge inside the QGP. The $r_D(T)$ is the Debye radius, which quantifies the specific extension of the color screening effect in the deconfined medium. The $r_D(T)$ strongly depends on the temperature of the system and is expected to decrease while the medium temperature increases. Therefore, the bound states which extend to a radius $r > r_D$ cannot survive in the QGP (Eq. 1.33). Moreover, the temperature dependence of the potential coefficients affects the effective radius of the bound state: according to Lattice QCD, at $T/T_c = 1.5$ the Colombian-like coefficient amounts to $\alpha \simeq 0.2$, consequently the J/ψ radius at these temperatures reaches within $0.5 \leq r_{J/\psi} \leq 1.3\text{fm}$ [100, 60]. A decrease of the Debye radius reduces the formation and the survival probability of quarkonia states in the QGP [47]. Since the binding energy of excited $q\bar{q}$ pairs ($\psi(2S), \Upsilon(2S), \Upsilon(3S)$) is lower than that of the quarkonium ground states ($J/\psi, \Upsilon(1S)$), a sequential production of the quarkonium resonances according to their binding energy is expected. This feature can be used to estimate the temperature of the medium according to the melting of specific bound states. Otherwise, the quarkonia dissociation, and a sequential decrease of the production yields, can be used as a thermometer of the QGP [2]. Some expressive values of dissociation temperature for the different quarkonia states are reported in Tab. 1.1

State	J/ψ	χ_c	$\psi(2S)$	$\Upsilon(1S)$	$\Upsilon(2S)$	$\Upsilon(3S)$
T_d/T_c	2.10	1.16	1.12	> 4.0	1.60	1.17

Table 1.1: Dissociation temperature in units of the critical temperature for different charmonia and bottomonia states. Numbers from [45].

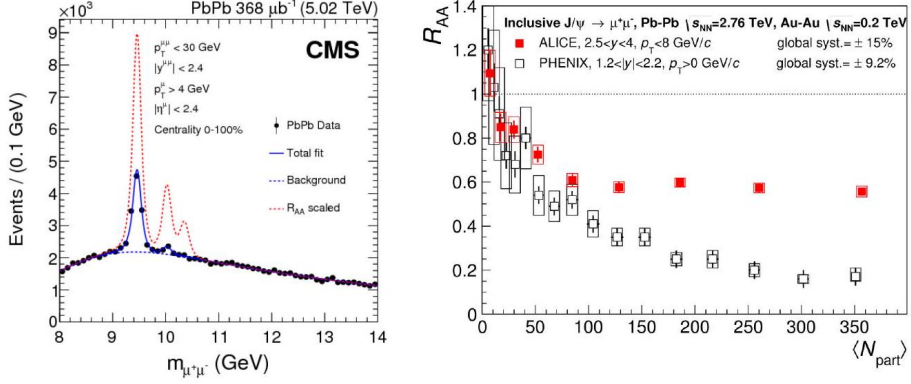


Figure 1.18: Left: invariant mass distribution of muon pairs in Pb – Pb collisions at $\sqrt{s_{NN}} = 5.02\text{TeV}$ measured by CMS. The measurements are shown as black points and fit with a solid blue line. The dotted red line is the same distribution scaled by the inverse of the nuclear modification factor. Figure from [19]. Right: nuclear modification factor as a function of $\langle N_{part} \rangle$ of J/ψ reconstructed in the $\mu^+\mu^-$ channel measured by the ALICE and PHENIX in Pb – Pb collisions at $\sqrt{s_{NN}} = 2.76\text{ TeV}$ and Au – Au collisions at $\sqrt{s_{NN}} = 200\text{ GeV}$ respectively. Figure from [20].

The J/ψ suppression was firstly observed at SPS in Pb – Pb collisions at $\sqrt{s_{NN}} = 17.4\text{ GeV}$ [101] and at RHIC in Au-Au collisions at $\sqrt{s_{NN}} = 200\text{ GeV}$ [102]. In the left panel of Fig. 1.18 the $\mu^+\mu^-$ invariant mass distribution in Pb – Pb collisions at $\sqrt{s_{NN}} = 5.02\text{ TeV}$ by CMS is shown [19]. The measured range is $8 < m_{\mu^+\mu^-} < 14\text{ GeV}/c^2$, to increase reconstructed signal for the bottomonia states $\Upsilon(1S)$, $\Upsilon(2S)$ and $\Upsilon(3S)$, with a mass of $m(\Upsilon(1S)) = 9.460\text{ GeV}/c^2$, $m(\Upsilon(2S)) = 10.023\text{ GeV}/c^2$ and $m(\Upsilon(3S)) = 10.355\text{ GeV}/c^2$, respectively [49]. From the comparison of the reconstructed distribution (black points and blue line) with the expected in the absence of QGP effects (red dotted line), the measurements show a decrease of the $\Upsilon(1S)$ yield of more than a factor $\times 2$, and the almost disappearance for the excited quarkonia states. These results confirm the hypothesis of a medium-induced color screening and the resulting sequential melting of quarkonia states. Consequently, quarkonium production is anticipated to be more suppressed with increasing temperature. This can be studied experimentally with the increase of the center-of-mass energy.

The nucleus-nucleus collisions at the LHC are, in principle, a perfect system to measure a decrease in quarkonia production with respect to pp collisions at the same energy and nucleus-nucleus collisions at low collision energy [47]. The comparison

of the nuclear modification factor for J/ψ mesons measured in Pb – Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV by the ALICE and in Au-Au collisions at $\sqrt{s_{NN}} = 200$ GeV measured by the PHENIX is shown in the right panel of Fig. 1.18. A larger J/ψ suppression is measured at RHIC, given a $R_{AA} \simeq 0.2$ for most central collisions versus a 3 times larger nuclear modification factor for Pb – Pb events at $\sqrt{s_{NN}} = 2.76$ TeV in correspondence of the same $\langle N_{part} \rangle$, differently from what was expected. These behaviors suggest a fundamental phenomenon that balances quarkonium suppression due to screening [47].

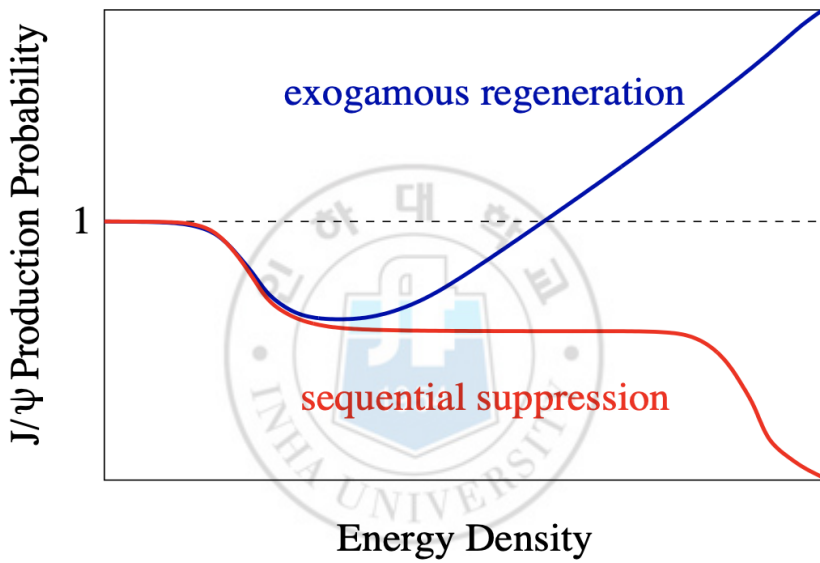


Figure 1.19: J/ψ sequential suppression compared with statistical recombination as a function of the system energy density. Figure from [21].

A production increase at LHC can be expected due to a more considerable quarkonium regeneration. According to the mechanism described in Ref. [103, 104, 105], the charmonia production comes from the statistical recombination of $c\bar{c}$ pair coming from unrelated hard scatterings. This contribution is anticipated to be more relevant at high center-of-mass energies due to the higher number of produced $q\bar{q}$ pairs, as shown in Fig. 1.19. The experimental measurements of R_{AA} as a function of p_T [20] show a nuclear modification factor about 4 times larger at LHC energies with respect to RHIC ones for $p_T(J/\psi) < 1$ GeV/ c , suggesting that the recombination is much more relevant for the charmonia production in the low p_T region [21].

Chapter 2

Physics of heavy-flavor production

Charm (c) and beauty (b) quarks are mainly produced in the initial hard scatterings, thanks to their high mass ($m_c \simeq 1.3 \text{ GeV}/c^2, m_b \simeq 4.2 \text{ GeV}/c^2$) [49]. These processes take place with a characteristic timescale $\Delta\tau \sim 1/Q$ of $\Delta\tau_c \sim 0.08 \text{ fm}/c$ and $\Delta\tau_b \sim 0.03 \text{ fm}/c^1$, which is smaller than the QGP formation time $\tau_{\text{QGP}} \sim 0.2 \div 1 \text{ fm}/c$ (see Sec. 1.3.2). Therefore, charm and beauty quark production is a useful tool to probe the QGP, and they survive through the medium expansion, interacting with the free partons in the plasma.

The formation of QGP can modify the hadronization mechanisms involving the heavy quarks. However, studying the hadronization process in pp collisions, where the QGP formation is not expected, is mandatory to understand the medium-induced effects thoroughly. These aspects are introduced in Sec. 2.1

In this Chapter, a brief introduction of the open heavy-flavor production in pp collisions is provided in Sec. 2.2 and the production quarkonium is also introduced in Sec. 2.3

2.1 Heavy-flavor production

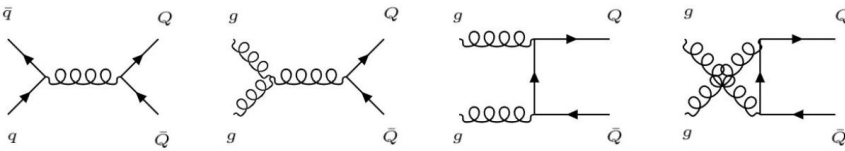


Figure 2.1: Feynman diagram of leading order quark-antiquark annihilation and gluon fusion for the production of $c\bar{c}$ and $b\bar{b}$. Figure from [22].

The pp collisions system is used to test the pQCD predictions, given that the large Q^2 required for the $c\bar{c}$ and $b\bar{b}$ pair production suggests a value for α_s which is significantly lower than unity (see Sec. 1.1). Moreover, measuring the open heavy-flavor hadron production in pp collisions provides the baseline for larger collision

systems. In such conditions, perturbative expansions in α_S can be engaged in the $q\bar{q}$ production cross section calculation. The main contributing processes at leading order ($\mathcal{O}(\alpha_S^2)$) are the quark-antiquark annihilation $\sigma_{q\bar{q} \rightarrow Q\bar{Q}}$ and the gluon fusion $\sigma_{gg \rightarrow Q\bar{Q}}$ [22], as shown in Fig. 2.1. In pQCD calculations the production cross section of heavy-flavor hadrons (H_Q) $d\sigma_{pp \rightarrow H_Q+X}$ is calculated by a factorization theorem, as following equation:

$$d\sigma_{pp \rightarrow H_Q+X}(\sqrt{s}) = \sum_{i,j=q,\bar{q},g} f_1(x_i, \mu_F^2) f_2(x_j, \mu_F^2) \otimes d\sigma_{ij \rightarrow h \rightarrow Q\bar{Q}}^h(\alpha_S(\mu_R^2), \mu_F^2, m_Q, x_i x_j s) \otimes D_Q^{H_Q}(z, \mu_F^2) \quad (2.1)$$

The convolutions of three terms in Eq. 2.1 are explained by the factorization theorem [106]. Considering two colliding protons, labeled as "1" and "2", the first term is $f_1(x_i, \mu_F^2) f_2(x_j, \mu_F^2)$, where $f_{1,2}$ are the parton distribution functions (PDFs) of the two colliding particles. They represent the probability for a given parton to carry a certain fraction of momentum of the original proton $x = p_{\text{parton}}/p_p$, called Bjorken- x . The second term $\sigma_{ij \rightarrow h \rightarrow Q\bar{Q}}^h$ is the cross section of the hard scattering process, where the heavy- $q\bar{q}$ pair is produced by the interaction of partons. The third term $D_Q^{H_Q}(z, \mu_F^2)$ is the fragmentation function, which is the probability for a heavy quark Q to fragment into a hadron H_Q carrying a fraction $z = p_{H_Q}/p_Q$ of the original quark momentum [106]. Heavy quarks and anti-quarks are separately fragmented into hadrons for heavy flavor production. The fragmentation of heavy quarks does not depend on the mechanism by which heavy quarks are produced [106]. For quarkonium production, typically, only 1-2% of the heavy quark anti-quark pairs form a bound quarkonium state instead of a pair of hadrons with open heavy flavor. The hadronization into the quarkonium bound states contains energy scales of the heavy quark momentum and the binding energy.

2.1.1 Single charm hadron in pp collisions

In Eq. 2.1, Both the PDF and the fragmentation functions depend on the virtuality (Q^2) of the process, linked to the factorization scale μ_F . PDFs are usually parametrized from measurements of deep-inelastic scattering processes [107, 108, 109] and the fragmentation functions parametrized from hadron production measurements in e^+e^- collisions [110, 111]. This approach is valid under the assumption of the universality of the fragmentation functions among the collision systems. The hard scattering cross section is a perturbative series on α_S , which depends on the energy scale, which is indicated by the renormalization scale μ_R defined for the

ultraviolet divergences renormalization in loop calculations [3]. Both the factorization and normalization scales are usually set to the same order of the momentum transfer of the hard process $\mu_F \sim \mu_R \sim \mu_0 \equiv \sqrt{m_Q^2 + p_{T,Q}^2}$. In addition, apart on the factorization scale μ_F and the bare value of the heavy mass m_Q , the cross section depends on the pp collision energy square $s = 4E_p^2 \simeq 4p_p^2$, which at the partonic level reflects in $x_i x_j s$ [106]. At present, the perturbative calculations are done up to the next-to-leading order $\mathcal{O}(\alpha_s^3)$ (NLO) and at a fixed order with the next-to-leading-log resummation in GM-VFNS [112, 113] and in FONLL [114, 115]. In particular, the next-to-leading-log corrections are more appropriate at high p_T , where $p_T \gg m$ given the quark mass m . Recently, calculation for the $c\bar{c}$ and $b\bar{b}$ production cross section are provided via next-to-next-to-leading-order (NNLO) calculations with QCD radiative corrections [116, 117, 118, 119, 120].

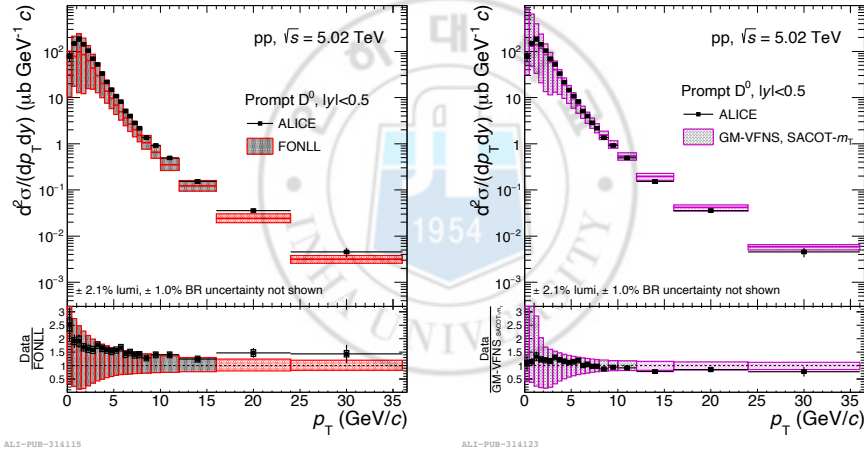


Figure 2.2: p_T -differential production cross section of D^0 mesons in pp collisions at $\sqrt{s} = 5.02$ TeV measured by the ALICE experiment compared with theoretical predictions from different pQCD calculations (details in the text). Figures from [23].

The factorization approach is tried by comparing model predictions and measurements of the open heavy-flavor hadron production. The results from the ALICE for the $D^0(c\bar{u})$ meson production cross section in pp collisions at $\sqrt{s} = 5.02$ TeV are shown in Fig. 2.2. The D^0 cross section is compared with FONLL and GM-VFNS calculations in Fig. 2.2. Different PDFs [107, 108] and fragmentation functions are considered in the two ways. The two calculations assume a different value for the charm quark mass, that is set to $m_c = 1.5 \text{ GeV}/c^2$ in FONLL and $m_c = 1.3 \text{ GeV}/c^2$

in GM-VFNS. The error bars around the theoretical calculation derive from the PDF and fragmentation function uncertainties and the scale variation of the mass of the charm quark. The two predictions explain the measured D^0 cross section well while systematically on the upper edge of the FONLL-based calculations. These results suggest that the factorization method works for the meson sector, as also observed by other LHC experiments [121, 122, 123]. However, the charm baryon results do not show the same conclusion. This behavior will be discussed in Sec. 2.2.

2.1.2 Multi-charm hadron in pp collisions

The measurement of the production yields of multi-charm baryons, which a mixture of uncorrelated charm quarks can only produce, would provide a new view on the production of heavy-flavor hadrons. Measurements of multi-charm hadrons, such as the Ξ_{cc}^+ (ccd), Ξ_{cc}^{++} (ccu), Ω_{cc}^+ (ccs), Ω_{ccc}^{+++} (ccd), and exotic states such as the T_{cc}^+ (ccud) [124, 125], would provide a clear view of hadron formation from QGP. The yields of multi-charm baryons relative to the number of produced charm quarks are predicted to be significantly enhanced in AA with respect to pp collisions. Enhancements are predicted by as much as a factor 102 for the recently observed Ξ_{cc}^+ baryon [126] and even by as much as a factor 1003 for the as yet undiscovered Ω_{ccc}^{+++} baryon. The observation and precise quantification of these effects will provide the study of deconfined matter's properties.

Fig. 2.3 shows the ratio of multi-charm baryons and the Λ_c yield as a function of the colliding system size, indicating the considerable sensitivity of these yields to the production mechanisms. For example, the SHM calculation assumes full equilibration of the charm quarks. Comparing the measured yields to these predictions would provide a compassionate measure of the degree of equilibration of charm quarks in the medium, including its system-size dependence. Ultimately, it must be possible to connect the equilibration to the transport properties of the QGP via kinetic transport modeling.

2.1.3 Quarkonium Production

Quarkonium production is an essential probe for studying both perturbative and non-perturbative QCD. The $q\bar{q}$ pair production and its evolution into quarkonium state affect both perturbative effects of short-distance/high-momentum and non-perturbative effects of long-distance/low-momentum. The factorization approach allows the separation of the perturbative and non-perturbative effects. The production of $q\bar{q}$ pair can be described similarly to the open heavy flavor production since the process requires large momentum transfer in the order of the heavy quark mass m_Q

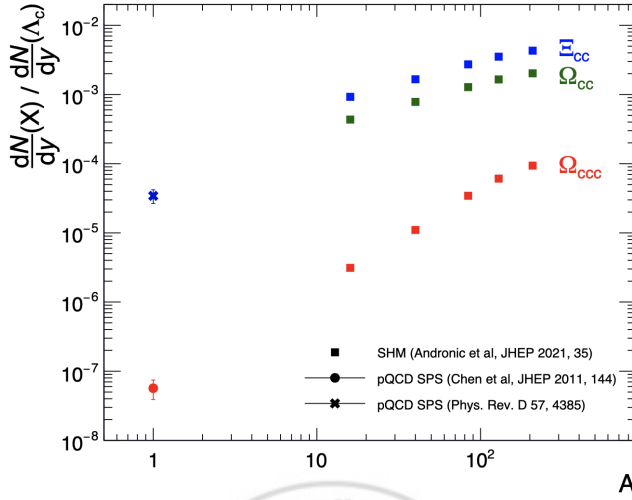


Figure 2.3: The wide span between the single parton scattering expectation and the statistical hadronization model prediction for the ratio of the multi-charm baryons and the Λ_c yield as a function of the colliding system size. Figures from [24].

or even larger scales. On the other hand, the $q\bar{q}$ pair evolution into the quarkonium state is a non-perturbative process requires over long distances, at the momentum scales of the heavy-quarks ($m_q v_q$) or smaller. Furthermore, their binding energy $m_q v_q^2$, where v_q is the velocity of the heavy quark in the bound state. The following scales describe the quarkonium states:

- Heavy-quark mass m_q : Treated perturbatively due to its hard scale, with a short-distance scale for annihilation processes.
- Momentum $m_q v_q$: Momentum of the components in the bound state and also the inverse of the size of the quarkonium state, usually treated non-perturbatively.
- Binding energy $m_q v_q^2$: Kinetic energy in the scale of the energy splittings between the radial excitations and the orbital angular-momentum excitations.

The spin splittings for a given radial and orbital angular-momentum excitations are in the order of $m_q v_q^4$, but the scale is negligible. The velocity of the heavy quark v_q decreases as the mass of heavy quark m_q increases. When the mass is enough, the velocity of the heavy quark becomes proportional to the running coupling constant $\alpha_s(m_q)$. Therefore, v_q decreases asymptotically proportional to $1/\ln m_q$, as in

Eq. 1.8. Also, When the mass is enough, the scales can be separated as $(m_q v_q^2)^2 \ll (m_q v_q)^2 \ll M_q^2$. Based on the quark potential model calculations, v_q^2 of the charmonium state is estimated to be about 0.3, for the bottomonium state is to be about 0.1 [127, 128]. If the running coupling constant is to be much smaller than 1, the effects at the momentum scale m_q and lower momentum scales such as $m_q v_q$ and $m_q v_q^2$ can be split. The assumption that $\alpha_s(m_q) \ll 1$ is well agreed with bottomonia where $\alpha_s(m_b) \simeq 0.18$, and also agreed with charmonia where $\alpha_s(m_c) \simeq 0.24$ [127]. As a result, it is possible to treat a quarkonium state as a non-relativistic system. The factorization involves the separation between the relativistic physics of annihilation for the heavy quark and anti-quark (for momenta $p \sim M_q$) and the non-relativistic quarkonium system for $p \sim M_q v_q$. The effective field theory separates different energy scales, non-relativistic quantum chromodynamics (NRQCD) [129, 130, 131]. The $q\bar{q}$ production requires a momentum transfer of the order M_q or larger, the inclusive cross section for the direct production of quarkonium H can be written as a sum of products of NRQCD matrix elements and short-distance cross section as below:

$$\sigma[H] = \sum_n \sigma_n(\Lambda) \langle \mathcal{O}_n^H(\Lambda) \rangle \quad (2.2)$$

where σ_n is the short-distance partonic cross section, which is convolved with the parton distributions of hadrons in the initial state. Λ is the ultraviolet cutoff from the effective theory [132, 133] and $\mathcal{O}_n^H$ are the expectation values of four-fermion operators in NRQCD [134].

The $q\bar{q}$ pair production is available in the color-singlet and color-octet states. Its spin state can be either singlet or triplet, and it is possible to have an orbital angular momentum. The color-singlet and color-octet four-fermion operators in Eq. 2.2 correspond to the evolution into a color-singlet quarkonium of a $q\bar{q}$ pair created at a short distance in a color-singlet state or a color-octet state, respectively. In the following, three different theoretical models are introduced to describe the non-perturbative dynamics in quarkonium production in Sec. 2.3.

2.2 Charm hadron production mechanisms

The string fragmentation in the Lund model [135] had been successfully described the hadron production in e^+e^- and e^-p collisions [110, 111]. However, the hadron production in hadronic collisions at the LHC energies can not be described by model prediction, especially for the production of charmed baryons.

In Fig. 2.4 the Λ_c^+/D^0 ratios in pp collisions at $\sqrt{s_{\text{NN}}} = 5.02$ [25] measured by

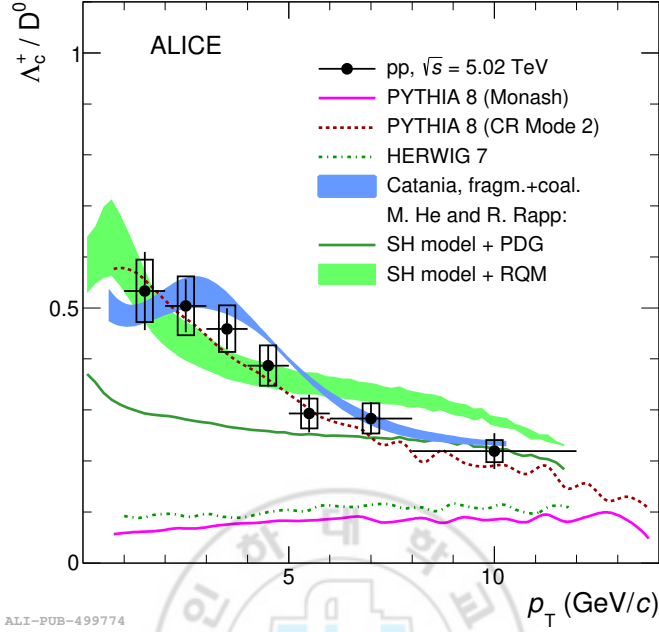


Figure 2.4: Λ_c^+ / D^0 ratio in pp collisions at $\sqrt{s_{NN}} = 5.02$ compared with predictions. Figure from [25]

the ALICE are shown. The measurements are largely underestimated by the model calculations of the PYTHIA8 Monash tune, which used the traditional fragmentation, which was measured in e^+e^- collisions. Especially, the measurements show the enhancement by about a factor of 5 at low p_T compared with model calculation. In order to understand this behavior, other model calculations need to be considered, and several models different from the Lund model have been suggested.

2.2.1 PYTHIA CR Mode 0, 2, 3

Traditionally, Monte Carlo generators use the leading-color (LC) approximation, which traces the color flow on an event-by-event basis. This approach considers that each quark is connected only to another colored parton in the event, similar to a leading-color QCD dipole [136], where the gluons carry a color and an anti-color charge, make connections with each other partons. Strings can be interpreted as gluons forming transverse "kinks" on strings whose endpoints are a quark and an antiquark by the Lund model [135]. This approach is challenged in hadronic collisions systems, such as pp collisions while showing a good agreement to describe the

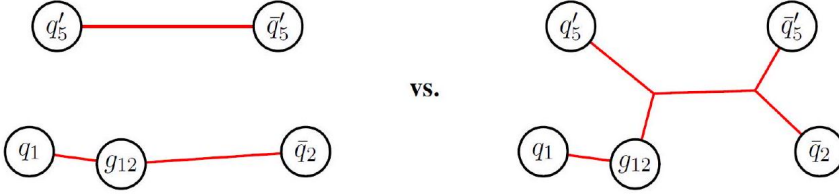


Figure 2.5: Schematic of a multi-parton state with a leading color string topology and with a junction-antijunction string topology, allowed by cyclically matching "2" and "5" indices. Figure from [26].

hadronization in e^+e^- collisions. In pp collisions, the initial state is sensitive to the strong force, and the partons inside protons carry a color charge, and colored beam leftovers remain after the collision. Moreover, multiple parton interactions (MPI) can occur, and they are necessary to explain the underlying event and soft-inclusive physics, especially in high-multiplicity events. The possibility of phase-space overlaps among particles produced in different MPI events is nonnegligible.

The model predictions [137, 138] already considered the possibility for the different MPI systems to be color-connected within a color chain to minimize the total color charge. However, the effect of color reconnection was estimated to be insignificant, without notable changes in the per-particle spectra in high-multiplicity events (many MPIs) with respect to low multiplicity events (low MPIs). In particular, this approach estimates the average p_T of charged particle for such events to be flat as a function of particle multiplicities. However, the experimental measurements [139, 140, 141] exclude this possibility, suggesting that need to introduce non-trivial color reconnection effects. The pure $SU(3)$ group structure of QCD already provides some sub-leading combinations among different partons that may create color singlet before the single parton fragmentation. For example, considering three uncorrelated color triplets (i.e., quarks), the Lie algebra of the $SU(3)$ group foresees:

$$3 \otimes 3 \otimes 3 = 10 \oplus 8 \oplus 8 \oplus 1 \quad (2.3)$$

In this equation, the decuplet corresponds to the LC topology, and one or two strings are added to connect to the corresponding beam remnant. This is done incoherently; no strings are extended among different MPI systems. The other multiplets describe additional sub-leading topologies that are not affected by fragmentation, and, in particular, one of them already produces a color singlet corresponding to a baryon. However, these additional combinations are not enough to describe the

baryon production in hadronic collision. Therefore, an additional color reconnection has been implemented beyond leading color approximation (CR-BLC).

This model [26] assumes a "simplified QCD" framework, in which parton combinations are ruled by 9 color indices, leading to 9 possible states for quarks and 72 for gluons, which allow recombinations among partons in different MPIs and from beam remnants. In this new framework, all the possible combinations among partons from the whole event are considered, enhancing the string topology involved in hadronization. These additional combinations are the so-called "junctions".

They are reported in the right of Fig. 2.5 and augment the set of possible string topologies that would be available with the LC structure. The possible topologies that can be realized in a specific event are governed by the minimization of the λ string-length measure, which corresponds to the area crossed by the string before the hadronization, i.e., the potential energy of the string. A necessary condition for junction connection is that its invariant mass must overcome a mass threshold m_0 corresponding to the number of possible recombinations. The parameter $C_j \equiv m_{0j}/m_0$ governs the amount of junctions, where to a high m_{0j} corresponds a low λ , resulting enhancement of junctions, $C_j > 1$.

Another component that affects the string topology is the causal connection among the strings themselves. If either string has already hadronized before the formation, there is no space-time region available for the reconnection among them. If two strings can reconnect, they must resolve between the string formation and its hadronization, considering the time dilation due to the relative boost. The causal contact between the two strings can be represented in terms of the following relation for the relative boost parameter γ :

$$\gamma \tau_{\text{form.}} < C_{\text{time}} \tau_{\text{had.}} \Rightarrow \frac{\gamma c}{m_{\text{string}} r_{\text{had.}}} < C_{\text{time}} \quad (2.4)$$

where $\tau_{\text{form.}} \sim 1/m_{\text{string}}$ is the proper formation time of a string and $\tau_{\text{had.}} = r_{\text{had.}}/c \sim 1/\Lambda_{\text{QCD}}$ is proper hadronization time, and C_{time} is a tunable parameter.

This model's different modes depending on the causal connection's requirement and strength. In Mode 0, no time-dilation condition is applied, but the amount of color reconnection is described only by the invariant-mass threshold m_0 . In Mode 2, the strictest condition is applied. The causal contact must be valid among all the dipoles involved in reconnection. In Mode 3, the constraint is more relaxed than Mode 2. Only the first reconnection between two dipoles is considered to satisfy Eq. 2.4, even if further recombinations occur later. The parameters of the model were tuned using experimental results: for example, the C_j parameter which is the number of junctions was tuned to reproduce the Λ/K_S^0 data in pp collisions at $\sqrt{s} = 7 \text{ TeV}$ by the CMS [142]. According to these model tunes based on the measurements, the

junction topology significantly enhances the production of open charmed baryons in pp collisions at the LHC energies, as shown in Fig. 2.4.

2.2.2 Statistical Hadronization Model with Relativistic Quark Model states

In the statistical hadronization model (SHM) [72], the hadron production is treated with the thermo-statistical weights. In SHM, the hadron production is governed only by the hadron mass at a given temperature of T_H . This approach has successfully described the light and strange hadron production in heavy-ion collisions and smaller systems, with the introduction of a strangeness suppression factor $\gamma_s < 1$ (see Sec. 1.4.5). However, the description of the baryon-to-meson ratio by SHM considering only the charm-hadron states reported in the PDG [49] cannot describe the Λ_c^+/D^0 ratio measured by the ALICE in pp collisions at $\sqrt{s} = 7\text{ TeV}$ [143], which shows enhancement with respect to the model calculations. The model in Ref. [46] provides a reconsidered framework that includes an augmented set of charm-baryon states expected by the Relativistic Quark Model (RQM [144]). They consist in 18 extra Λ_c^+ states, 42 Σ_c 's, 62 Ξ_c 's and 34 Ω_c 's up to a mass of $3.5\text{ GeV}/c^2$. Thermal densities, which considered all hadron states, are governed by supervised by specific weights that depend on the hadron mass m_i and their spin degeneracy $d_i = 2J + 1$:

$$n_i = \frac{d_i}{2\pi^2} m_i^2 T_H K_2 \left(\frac{m_i}{T_H} \right) \quad (2.5)$$

where K_2 is the second order modified Bessel function, and T_H is the hadronization temperature.

The upper limit of T_H is assumed to be 170 MeV , according to the results in Ref. [145]. The thermal densities in Eq. 2.5 are used to designate the proper normalization to the charm-quark fragmentation functions, which are extrapolated from those assumed in the FONLL framework, employed to compute the charm p_T spectrum produced in pp collisions. All the additional excited charmed baryons are forced to decay into a ground state. This process does not change the abundance of the ground state D mesons. However, it significantly enhances the ground state charmed baryons, as reported in Tab. 2.1.

2.2.3 Catania coalescence model

Recent measurements in pp collisions at the LHC [85, 146, 121, 147], show a phenomenology similar to that observed in AA collisions, such as strangeness

$n_i (\times 10^{-4} \text{fm}^{-3})$	D^0	D^+	D^{*+}	D_s^+	Λ_c^+	$\Xi_c^{0,+}$	Ω_c^0
PDG ($T_H = 170 \text{MeV}$)	1.161	0.5098	0.5010	0.3165	0.3310	0.0874	0.0064
RQM ($T_H = 170 \text{MeV}$)	1.161	0.5098	0.5010	0.3165	0.6613	0.1173	0.0144

Table 2.1: Thermal densities estimated by the SHM model for prompt ground state charmed hadrons at the hadronization temperature $T_H = 170 \text{MeV}$ including only hadrons listed in the PDG and also adding the ones expected by the RQM. Numbers from [46].

and baryon-to-meson enhancement, especially when particle yields are studied as a function of event multiplicity. Different theoretical studies [148, 149, 150, 151] considering hydrodynamical and transport-based calculations show reasonable description of p_T spectra and two-particle correlations, pointing to a possible formation of a hot QCD matter with a lifetime $\tau \approx 2 \text{fm}/c$. Based on the hydrodynamic approach, the Catania model assumes the presence of a thermalized system of gluons and u, d, s quarks and antiquarks already in pp collisions, characterized by a lifetime $\tau = 2.5 \text{fm}/c$, a transverse radius $R = 2 \text{fm}$ and a temperature $T = 165 \text{MeV}$, according to the hydrodynamical simulations in Ref. [148]. Since the charm thermalization time is larger, no thermalized charm quarks are expected in bulk. The presence of a deconfined medium is crucial for the Catania model, where hybrid hadronization via fragmentation and coalescence is assumed for the charm quark. The probability of coalescence of charm quark with a light parton in bulk is complementary to the fragmentation probability, $P_{\text{fragm.}} = 1 - P_{\text{coal.}}$. In this model, the coalescence is described with the Wigner formalism, according to which the p_T spectrum of a hadron produced by the charm coalescence can be described as:

$$\begin{aligned} \frac{d^3 N_H^{\text{coal.}}}{dy d^2 p_T} &= g_H \int \prod_{i=1}^{N_q} \frac{d^3 p_i}{(2\pi)^3 E_i} p_i \cdot d\sigma_i f_{q_i}(x_i, p_i) \\ &\times f_H(x_1, \dots, x_{N_q}, p_1, \dots, p_{N_q}) \delta^{(2)}\left(p_T - \sum_{i=1}^n p_{T,i}\right). \end{aligned} \quad (2.6)$$

where g_H indicates the statistical degeneracy factor to form a colorless hadron from quarks and antiquarks with spin 1/2 and corresponds to 1/36 for mesons and 1/108 for baryons. $d\sigma_i$ is an element of a space-like hypersurface, and $f_{q,i}$ is the (anti)quark phase-space distribution for the i -th (anti)quark. $f_H(x_1, \dots, x_{N_q}, p_1, \dots, p_{N_q})$ corresponds to the Wigner function, which describes the probability for a hadron to be produced from quarks with given phase-space coordinates.

The Eq. 2.6 represents the production via coalescence of mesons and baryons

in case of $N_q = 2$ and 3, respectively. The Wigner function is assumed to follow a Gaussian shape in space and momentum:

$$f_H(x_1, \dots, x_{N_q}, p_1, \dots, p_{N_q}) = \prod_{i=1}^{N_q-1} A_W \exp\left(-\frac{x_{r_i}^2}{\sigma_{r_i}^2} - p_{r_i}^2 \sigma_{r_i}^2\right) \quad (2.7)$$

where A_W is a normalization factor fixed to guarantee that for $p \rightarrow 0$, a charm quark indeed hadronises via coalescence. The satisfaction $\lim_{p \rightarrow 0} P_{\text{coal.}}^{\text{tot.}} = 1$ is required for total coalescence probability. The Wigner function depends on x_{r_i} and p_{r_i} , which are the relative phase-space coordinates among the valence quarks in the hadron and σ_{r_i} are covariant widths, related to the root-mean-square charge radius of the hadron: $\langle r^2 \rangle_{ch} = \sum_{i=1}^N Q_i \langle (x_i - X_{\text{CM}})^2 \rangle$, with $N = 2$ and 3 for mesons and baryons respectively. The hadron radii considered in this model are taken from the quark model [152, 153].

The p_T spectrum of charm hadron produced from the charm fragmentation can be written as follows:

$$\frac{d^3 N_H^{\text{fragm.}}}{dy d^2 p_T} = \Sigma \int dz \frac{d N^{\text{fragm.}}}{dy d^2 p_T} \frac{D_{H/c}(z, Q^2)}{z^2} \quad (2.8)$$

where $D_{H/c}(z, Q^2)$ corresponds to the fragmentation function, which is the probability of the charm quark to produce a hadron H with a fraction z of the original charm momentum, given $Q^2 = (p_{\text{had.}}/2z)^2$ the momentum scale of the fragmentation process.

These functions correspond to the Ref. [154], taken from Ref. [155] for the charm quark. The interplay between fragmentation and coalescence relies on the charm p_T .

In Fig. 2.6, the coalescence probability of the charm quark to the various hadron species is shown as a function of the charm p_T . The total probability reaches unity for $p \rightarrow 0$ as required in the definition of A_W in Eq. 2.7 and rapidly decreases with increasing charm p_T , indicating that at low p_T the charm quark only coalesces, while at high p_T the fragmentation contribution becomes dominant. It is noticeable that for $p_T < 2 \text{ GeV}/c$, the coalescence probabilities into Λ_c^+ and Ξ_c are compatible and are higher than the probability into a D^0 . The model calculation explain the enhancement of the Λ_c^+/D^0 , Ξ_c/D^0 and $\text{BR} \times \Omega_c^0/D^0$ ratios. In particular, the model calculations show the result $P_{\text{fragm.}} > P_{\text{coal.}}$ for the charm mesons and $P_{\text{fragm.}} < P_{\text{coal.}}$ for baryons, where a hadron p_T is lower than $10 \text{ GeV}/c$.

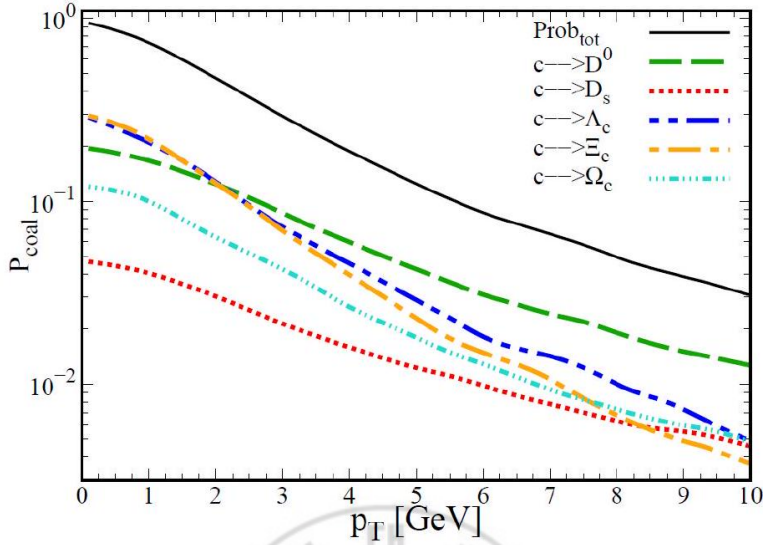


Figure 2.6: Charm coalescence probability as a function of the charm p_T in pp collisions at the LHC. The colored lines correspond to the probability of the charm quark coalescence for a given hadron species, while the solid black line is the total coalescence probability. Figure from [27].

2.2.4 Quark (re-)Combination Mechanism

In this model, the production of low- p_T charm hadrons originated as the coalescence of a perturbatively created charm quark with equal-velocity light quark deriving from fragmentation. The charm mesons and baryons are produced with a characteristic momentum $p_H = p_c + p_{\bar{q},qq}$ from a charm quark (p_c) that picks up a co-moving light antiquark ($p_{\bar{q}}$) or two co-moving light quarks (p_{qq}) respectively. The momentum distribution of a single-charm meson $M_{c\bar{l}}$ and baryon $B_{cll'}$, with l and l' indicating two light-flavour quarks among u, d and s, is written as:

$$f_{M_{c\bar{l}}}(p) = \int dp_1 dp_2 f_{ci}(p_1, p_2) \mathcal{R}_{M_{c\bar{l}}}(p_1, p_2; p) \quad (2.9)$$

$$f_{B_{cll'}}(p) = \int dp_1 dp_2 dp_3 f_{c ll'}(p_1, p_2, p_3) \mathcal{R}_{B_{cll'}}(p_1, p_2, p_3; p) \quad (2.10)$$

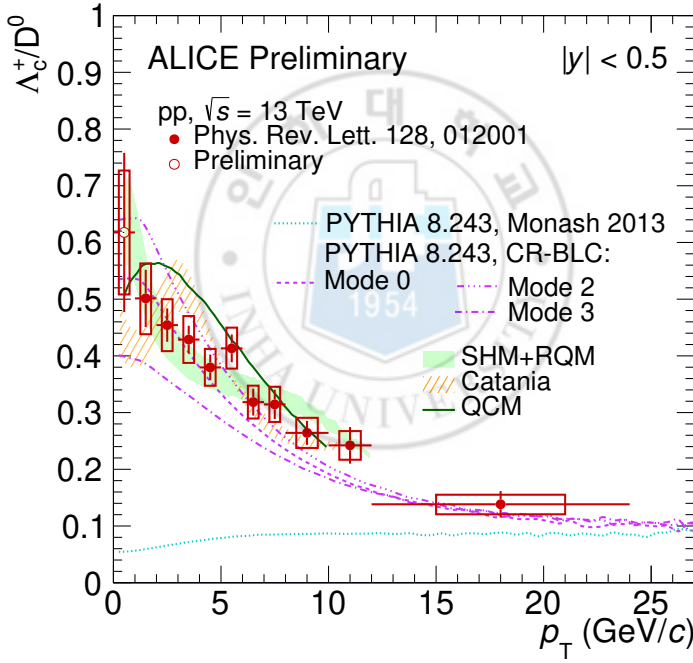
where $f_{ci}(p_1, p_2)$ is the joint momentum distribution for charm quark and light antiquark and $\mathcal{R}_{M_{c\bar{l}}}(p_1, p_2; p)$ corresponds to the probability density for the c and $\bar{l}$ with momenta p_1 and p_2 respectively to be combined into a meson $M_{c\bar{l}}$ and baryon $B_{cll'}$ with momentum p .

The momentum integrated yields for the charm meson $M_{c\bar{l}}$ and baryon $N_{B_{cll'}}$ are:

$$N_{M_{c\bar{l}}} = \kappa_{M_{c\bar{l}}} \int dp f_c(x_1 p) f_{\bar{l}}(x_2 p) = N_c N_{\bar{l}} \frac{\kappa_{M_{c\bar{l}}}}{A_{M_{c\bar{l}}}} = N_c N_{\bar{l}} \mathcal{R}_{c\bar{l} \rightarrow M_{c\bar{l}}} \quad (2.11)$$

$$N_{B_{cll'}} = N_c N_l N_{l'} \frac{\kappa_{B_{cll'}}}{A_{B_{cll'}}} = N_c N_l N_{l'} \mathcal{R}_{cll' \rightarrow B_{cll'}} \quad (2.12)$$

where $\mathcal{R}_{c\bar{l} \rightarrow M_{c\bar{l}}}$ and $\mathcal{R}_{cll' \rightarrow B_{cll'}}$ corresponds to the momentum-integrated probability for a c and $\bar{l}$ quarks to produce a meson $M_{c\bar{l}}$.



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Figure 2.7: Λ_c^+/D^0 ratio in pp collisions at $\sqrt{s} = 13$ TeV measured by the ALICE compared with the prediction of the QCM model (green line). Figures from ALICE Figure Repository.

Given these inputs, the QCM model provides predictions for the production of pseudo-scalar $J^P = 0^-$ (D^+ , D^0 , and D_s^+), and vector $J^P = 1^-$ (D^{*+} , D^{*0} , and D_s^{*+}) D mesons, as well as the triplet (Λ_c^+ , Ξ_c^+), and Ξ_c^0) with $J^P = (1/2)^+$, sextet (Σ_c^0 ,

$\Sigma_c^+, \Sigma_c^{++}, \Xi_c^{*0}, \Xi_c^{*+}$, and Ω_c^0) with $J^P = (1/2)^+$, and sextet ($\Sigma_c^{*0}, \Sigma_c^{*+}, \Sigma_c^{*++}, \Xi_c^{*0}, \Xi_c^{*+}$, and Ω_c^{*0}) with $J^P = (3/2)^+$ charm baryons.

In Fig. 2.9 the Λ_c^+/D^0 ratio measured by the ALICE in pp collisions at $\sqrt{s} = 13\text{TeV}$ [156] is compared with the prediction from the QCM model as a green line. The QCM describes both the magnitude and the p_T dependence within the uncertainties.

2.3 Quarkonium production mechanisms

The quarkonium production in the high-energy collision can be described by several sophisticated factors, including the model calculations. The following section introduces model calculations that describe the quarkonium production mechanisms, such as Color Evaporation Model (CEM), Color Single Model (CSM), and Non-relativistic QCD (NRQCD).

2.3.1 Color-Singlet Model

The color-singlet model (CSM) assumes that the quarkonium state is produced through the partonic processes producing a $q\bar{q}$ pair in a color-singlet state. The $q\bar{q}$ state has a colored state, and one hard gluon carries away the color to each quark. The color and spin of the $q\bar{q}$ state do not change from production to hadronization. Therefore, $q\bar{q}$ pair should have the same angular-momentum quantum numbers as the quarkonium state. The CSM can be obtained in Eq. 2.2 via the Non-relativistic QCD factorization, excluding all color-octet terms. The $q\bar{q}$ state production in CSM-corresponds to the color-singlet $q\bar{q}$ wave function and its derivatives, considered at zero $q\bar{q}$ splits determined by the annihilation decay rates. Thus, after the wave function at the origin is extracted from the measurements, the CSM has no free parameters and gives the normalized calculations for the production cross sections.

The CSM was successfully applied to the prediction of the quarkonium production rates at relatively low energies [157], e.g., the η_c and χ_c production through two-gluon fusion [158, 159, 160, 161] and the J/ψ and η_c production in B-meson decays [162, 163, 164]. However, the measurements from the Tevatron showed that the CSM significantly underestimated the production for the prompt charmonium production in pp collisions by more than a factor of 10. It turns out that considering only the leading order term of α_s and the v expansion of NRQCD makes the underestimation of the measurements in pp collisions. So the next-to-leading order (NLO) and next-to-next-to-leading order (NNLO) have to be considered, which significantly increases the quarkonium production [165, 28]. However, it is unclear whether the perturbative expansion in α_s is convergent or not, given the extensive

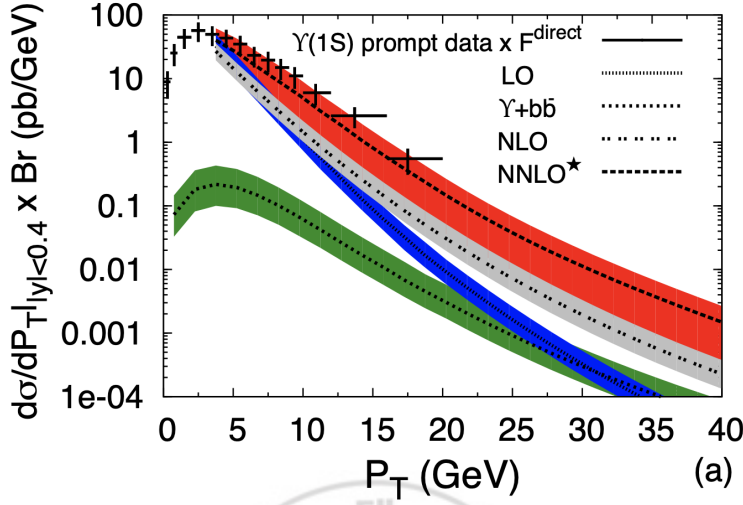


Figure 2.8: Model calculations for $\Upsilon(1S) + X$ and at LO, associated production, full NLO, and NNLO* compared with the direct yield at $\sqrt{s} = 1.8$ TeV measured by the CDF collaboration. The band represents the uncertainties of model calculations from the various parametrizations. Figure from [28].

corrections at NLO and NNLO. Furthermore, the CSM is unreliable in the production and decay of P -wave and higher-orbital-angular-momentum quarkonium states because it leads to uncanceled infrared divergences.

2.3.2 Color-Evaporation Model

The color-evaporation model (CEM) [166, 167, 47] suggests primary $q\bar{q}$ pair production with invariant mass from the quark threshold, $2m_q$, up to the open heavy flavor meson threshold. Afterward, the $q\bar{q}$ pair balances into the quarkonium state by emitting non-perturbative soft gluons. Thus, the color (also the spin) of the $q\bar{q}$ pair does not give constraints to the quarkonium formation. In CEM, the momentum of the quarkonium state is assumed to be the same as the momentum of the $q\bar{q}$ pair. The cross section of the quarkonium states H can be represented by the integration of the perturbative cross section for inclusive $q\bar{q}$ production in the CEM constraints mass range and formulated as:

$$\frac{d\sigma_H(P)}{d^3P} = F_H \int_{2m_q}^{2m_H} dm_{q\bar{q}} \frac{d\sigma_{q\bar{q}}(m_{q\bar{q}}, P)}{dm_{q\bar{q}} d^3P} \quad (2.13)$$

	J/ψ	$\psi(2S)$	χ_{c1}	χ_{c2}	$\Upsilon(1S)$	$\Upsilon(2S)$	$\Upsilon(3S)$	$\chi_b(1P)$	$\chi_b(2P)$
R_H	0.62	0.14	0.60	0.99	0.52	0.33	0.20	1.08	0.84

Table 2.2: Ratios of direct (prompt) to inclusive quarkonium production, $R_H = \sigma_{dir}^H / \sigma_{inc}^{H'}$. Numbers from [47].

where F_H is the fraction of the $q\bar{q}$ pair states evolving into the quarkonium state H .

The fraction F_H is universal and independent of the kinematics and spin, as well as the production process of the quarkonium state. Therefore, the quarkonium production ratios should be constant in all ranges of kinematic variables. In the case of charmonia, the fraction F_H of J/ψ has been measured in forward, and mid rapidity region, including all feed-down decays from the excited states such as χ_c and $\psi(2S)$. The measurement and branching ratios can obtain the fraction of the direct (prompt) quarkonium production compared to the inclusive cross section. The ratios of the direct (prompt) to inclusive cross section $R_H = \sigma_{dir}^H / \sigma_{inc}^{H'}$, where H' refers to the respective quarkonium ground state, are summarized in Table. 2.2.

The CEM has one parameter per quarkonium state, while the color-singlet model effectively has no free parameter. The simple CEM can successfully describe the J/ψ production [168]. However, the measurements show that the ratio of $\psi(2S)$ and J/ψ varies as a function of p_T [169, 170] can not be described the CEM prediction. In the recently modified CEM, the invariant mass of the $q\bar{q}$ pair is constrained to be larger than the mass of the formed quarkonium [171]. A momentum transition between the $q\bar{q}$ pair and quarkonium state has also been considered. The model calculations show better agreement to the charmonium production and the ratio of $\psi(2S)$ and J/ψ yield than the traditional CEM. Also, the polarization of the ground state of charmonia and bottomonia has been calculated, for the first time, in the color evaporation approach [171].

2.3.3 Non-Relativistic QCD

The relative small velocity of the quarkonium states, as discussed in Sec. 2.1.3, authorizes describing the quarkonium production with QCD in the non-relativistic limit based on an effective field theory, called NRQCD [131]. The NRQCD splits the short distances that are similar or smaller than the heavy-quark mass (m_q) from the universal long-distance matrix elements (LDMEs). Perturbative QCD can calculate the short-distance coefficients in the strong-coupling constant α_s . The velocity-scaling rules describe the long-distance effects in the limit ($v_q \ll 1$, which are used to determine the magnitude of the NRQCD matrix elements. In this way, the

NRQCD framework is established as double expansions of α_s and v_q . The cross section of a quarkonium state H can be described as the sum of the products of the short-distance coefficients and the LDMEs, similar to the characterization at the beginning of Sec. 2.1.3:

$$\sigma(H) = \sum_n \frac{F_n(\Lambda)}{m_q^{d_n-4}} \langle 0 | \mathcal{O}_n^H(\Lambda) | 0 \rangle \quad (2.14)$$

where Λ is the ultraviolet cutoff, F_n is the short-distance coefficient, and $\mathcal{O}_n^H$ is the non-perturbative operator with the mass dimensions of d_n . $F_n(\Lambda)$ is the partonic cross section that forms a $q\bar{q}$ pair in a color singlet or octet state.

The coefficients are obtained by matching the square of the production amplitude in NRQCD to the full QCD calculations done perturbatively. On the other hand, the operators in Eq 2.14 convert the produced $q\bar{q}$ pair into an intermediate state, composed of quarkonia and anything, and later annihilate the $q\bar{q}$ pair. The probability of forming the quarkonium plus anything from a $q\bar{q}$ pair is cited as the elements of the vacuum matrix. For $p_T \geq 2m_q$, NRQCD shows good agreement with the quarkonium production within uncertainties.

The production mechanism has not been fully understood, even considering the NRQCD descriptions for quarkonium production [172]. The expansion with the strong-coupling constant α_s and the velocity for the cross section of quarkonium states have been applied to the factorization of the hard and soft contributions and separation of different color and spin states. The elements of the long-distance matrix are weighted by the contributions from each color and spin state and then obtained from the fit to the measurement above a p_T threshold. In contrast to the universal prediction of the LDMEs, the elements depend on the collision system. Also, the distribution of the χ_c meson, measured in LHCb [29] can not be described by the model calculations from NRQCD [173].

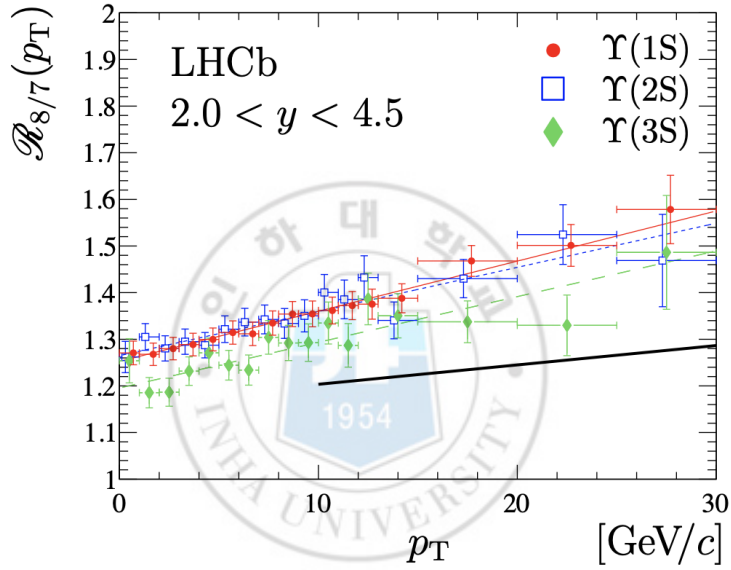


Figure 2.9: Ratios of the differential cross-sections $\sqrt{s} = 8$ and 7 TeV for three Upsilon states. The fit results with a linear function are shown with straight thin red solid, blue dotted, and green dashed lines. The next-to-leading order NRQCD theory predictions are shown in the same plot as a thick line. Figure from [29].

Chapter 3

Experimental setup

The ALICE (A Large Ion Collider Experiment) is one of the four major experiments at the Large Hadron Collider (LHC) of the CERN laboratory in Geneva, Switzerland. ALICE is dedicated to heavy-ion physics and the characterization of the Quark-Gluon Plasma. The Ξ_c^0 production in pp collisions at $\sqrt{s} = 13$ TeV, which are contained in this thesis, had been done using the data collected by ALICE detector. In this chapter, the experimental setup will be discussed installed in ALICE. The main features of LHC will be introduced in Sec. 3.1, with some parameters that are relevant to the proton-proton beam. The ALICE apparatus used to data collecting will be introduced in Sec. 3.2. The details of vertex and track reconstruction and the particle identification will be described in Sec. 3.4 and Sec. 3.5, respectively.

3.1 The Large Hadron Collider

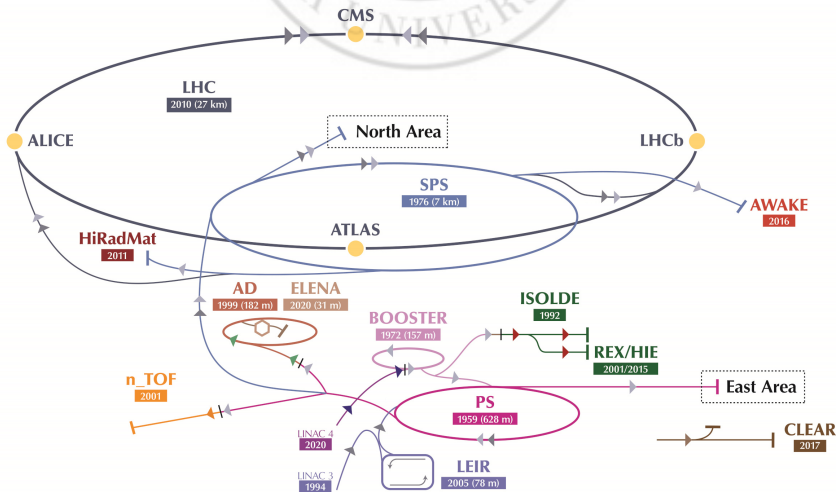


Figure 3.1: CERN accelerator complex [30].

The Large Hadron Collider (LHC) is the world's largest and highest energetic particle accelerator at the CERN laboratory located on the French-Swiss border. The collider is contained in a circular tunnel (27 km) and has two parallel beam pipes from 50 to 175 m underground. LHC is an acceleration complex composed of sub-accelerators of a sequential acceleration (Fig. 3.1). The accelerator is proposed to accelerate protons up to $\sqrt{s} = 14$ TeV, and heavy ions up to $\sqrt{s_{NN}} = 5.5$ TeV.

In high-energy physics, luminosity is a parameter expressing the interaction rate normalized to the intensity of the interacting fluxes of particles in the collider. The luminosity L can be expressed as:

$$L = R/\sigma \quad (3.1)$$

where R is the collision rate, and σ is the cross section for the beam collisions. The luminosity can be measured by the experimental technique, Van der Meer scan [174]. The luminosity can be defined with the Van der Meer scan as:

$$L = \frac{N_b^2 n_b f_{rev} \gamma F}{4\pi \epsilon_n \beta^*} \quad (3.2)$$

where N_b is the number of particles per bunch, n_b is the number of bunches per beam, f_{rev} is the revolution frequency, γ is the Lorentz factor, ϵ_n is the transverse beam emittance, β^* is the beta function at the interaction point, and F is a correction factor.

The targeting luminosities are different for each experiment by their experimental aim. The integrated luminosities of ALICE during the Run2 period, which is pp collisions at $\sqrt{s} = 13$ TeV, are shown in Fig. 3.2

3.2 ALICE: A Large Ion Collider Experiment

The ALICE experiment [31] is designed and optimized for the heavy ion experiment. The main goal of the ALICE experiment is to study the properties of the QGP and hadronization mechanism to search for the origin of matter. Therefore, ALICE needs to have a great performance of particle tracking and identification of all the particles produced in the heavy ion collisions. In this section, only specific detectors which are used for the analysis in this thesis will be described.

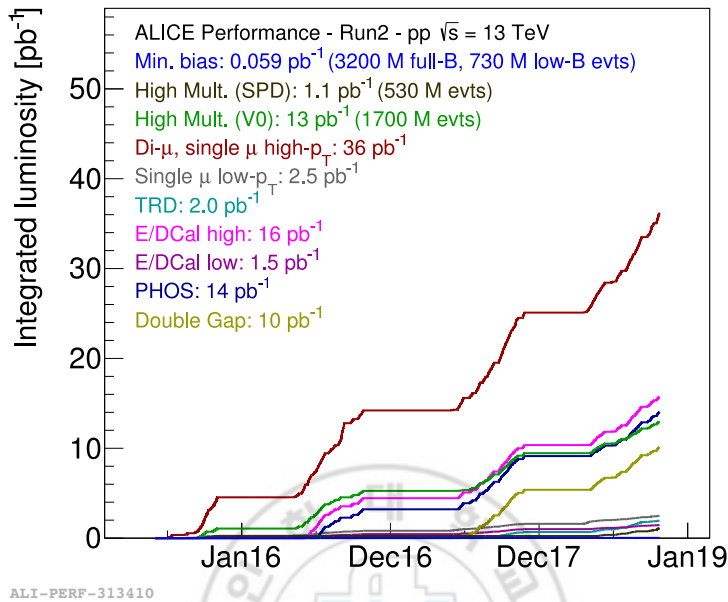


Figure 3.2: Integrated luminosity for various triggers in pp collisions at $\sqrt{s} = 13$ TeV in Run 2 period. Figure from ALICE figure repository.

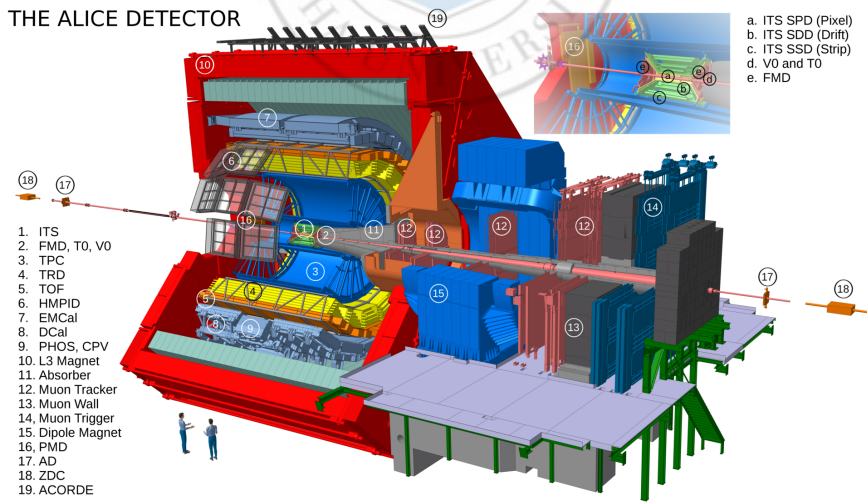


Figure 3.3: The ALICE apparatus for Run1 and Run2 with description. Figure from ALICE figure repository.

3.2.1 Inner Tracking System

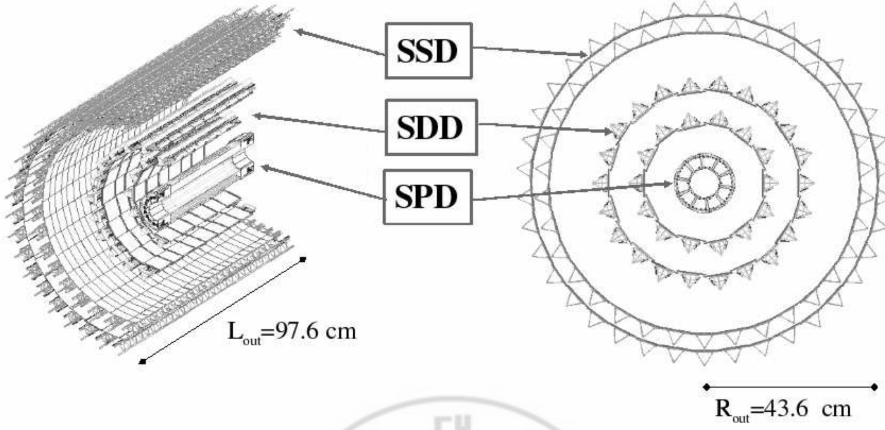


Figure 3.4: ITS schematic layout. Figure from [31].

The Inner Tracking System (ITS) [31] consists of six silicon layers located in the innermost of ALICE detector. The first two layers of the ITS are Silicon Pixel Detectors (SPD), located at 3.9 and 7.6 cm from the interaction point. They can cover the pseudorapidity range $|\eta| < 1.4$ and the full azimuth around the interaction point. The SPD is used to reconstruct SPD tracklet, which is a short track segment that is extended with two hits on each layer linearly to the interaction point of the two beams. The next two layers of ITS are Silicon Drift Detector (SDD), and the outermost layers of ITS are Silicon Strip Detector (SSD), with the outermost layer having a radius $r = 43.6$ cm.

The main task of the ITS is the primary and secondary vertices reconstruction, with a resolution better than $100 \mu\text{m}$. ITS provides the powerful tracking performance even for low p_T particles down to $p_T = 80 \text{ MeV}/c$, thanks to the close space points of the silicon layers.

3.2.2 Time Projection Chamber

The Time Projection Chamber (TPC) [175] is the main detector for track reconstruction and particle identification via energy loss of particles in TPC. TPC provides excellent momentum and spatial resolution for charged particles, and it is located just outside the ITS. The radial size of the TPC is about $85 < r < 250$ cm, and the longitudinal size is $-250 < z < 250$ cm. TPC can cover the pseudorapidity

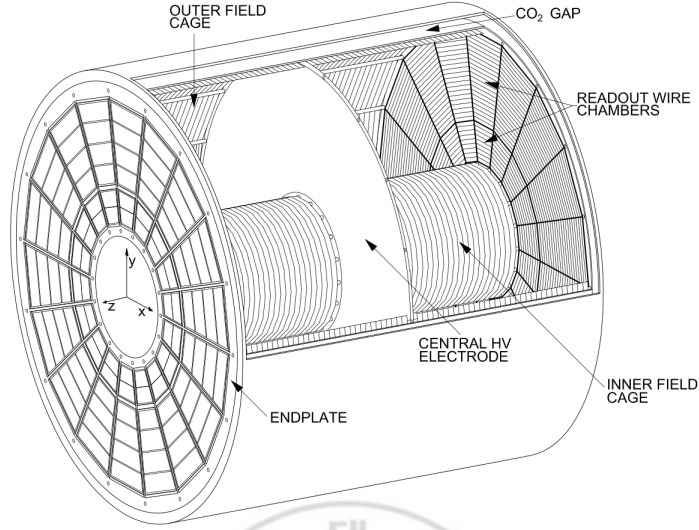


Figure 3.5: TPC schematic layout. Figure from [31].

range $|\eta| < 0.9$, and the full azimuth around the interaction point. TPC consists of a 90m^3 cylindrical drift chamber filled with a mixture of active gas (Ar, Ne, CO_2 , N_2).

The initial gas in the central barrel was a mixture with 90 % of Ne and 10 % of CO_2 , with the initialization of N_2 . This mixture was changed to Ar instead of Ne in 2016 and 2018, to reduce the space-charge distortion in TPC, which occurred by the charged particles.

3.2.3 Time-Of-Flight

The Time-Of-Flight detector (TOF) [175] is a large area detector located around the TRD with the same pseudorapidity coverage of ITS, TPC, and TRD ($|\eta| < 0.9$). TOF can identify the charged particles in the intermediate momentum range, up to $2.5\text{ GeV}/c$ ($4\text{ GeV}/c$). In the intermediate momentum range, pions (protons) and kaons separation is performed with more than 3 times resolution of time-of-flight. The TOF consists of Multi gap Resistive Plate Chambers (MRPC) which placed in an area of about 160m^2 , up and downside of the cylinder. The 10 layers double stack strip detectors are composed of multi gap with a time resolution of about 40ps. The time of flight (t_{flight}) for each particle can be derived using the time of particle hits the detector (t_{hit}) and the start time of the event, following this

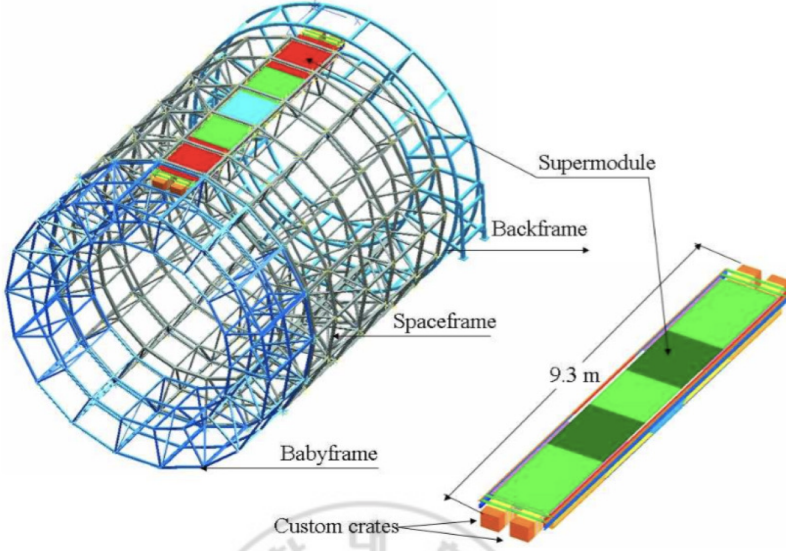


Figure 3.6: TOF schematic layout. Figure from [32].

relation, t_0 ($t_{flight} = t_{hit} - t_0$). The start time t_0 is determined by the T0 detector 3.2.4 or using the particle arrival times at the TOF, as described in [32]. The overall TOF resolution is around 100 ps in pp collisions, having uncertainty on the determination of the event time zero.

3.2.4 T0

The T0 detector [175] is two arrays of Cherenkov detectors (T0A, T0C), covering the full azimuth and the pseudorapidity range ($4.61 < \eta < 4.92$ and $-3.28 < \eta < -2.97$). T0 is used for the determination of the start time (t_0) for the TOF.

3.2.5 V0

The V0 detector [175] is composed of two arrays (V0A, V0C) of 32 scintillating counters. The V0A (V0C) is located at a distance of 329 cm (-88 cm) away from the nominal interaction point ($z = 0$), and they cover the pseudorapidity range $2.8 < \eta < 5.1$ ($-3.7 < \eta < -1.7$) and the full azimuth. V0 is used for event triggering and the determination of the Multiplicity (Centrality) classes of the pp (Pb-Pb) collisions with the sum of the signals deposited in the V0A and V0C.

3.2.6 The trigger and data acquisition

ALICE Trigger is performed by the Central Trigger Processor (CTP) [176], based on detector signals and information about the LHC bunch-filling scheme. A 3-level trigger system is deployed for triggering. The Level 0 (L0) trigger decision is done in $\sim 0.9 \mu\text{s}$ after the collision using fast detectors, such as V0, T0, EMCal, PHOS, and MTR. The Level 1 (L2) trigger decision is done after 260 LHC clock cycles ($\sim 6.5 \mu\text{s}$) for the event passing L0 trigger selection, due to slow computation time on the TRD and EMCal, and propagation times on ZDC. The Level 2 (L2) trigger decision is done after $\sim 100 \mu\text{s}$, due to the drift time of the TPC. All the events passing L2 selection are transported to the Data Acquisition (DAQ) machines and the High-Level Trigger (HLT), which are pure software triggers.

The minimum bias triggers (MB_{AND}) and a high-multiplicity (HM) trigger are used for all pp and Pb–Pb data taking on Run 2 period.

3.3 The ALICE offline framework

The ROOT [177] scientific software, which handles huge data processing, statistical analysis, data visualization, and storage for high-energy physics experiments, is the core of the ALICE offline software, known as AliRoot cite Aliroot. Although ROOT is based primarily on C++, it also connects with other languages like Python [178] and R [179]. The simulation, alignment, calibration, reconstruction, and analysis of the experimental data collected by ALICE are proposed by the AliRoot framework.

With the help of Monte Carlo generators, such as PYTHIA [180], and HERWIG [181] for pp collisions, simulations anchored to experimental data are generated. Each subdetector's detector geometry details are included in AliRoot and are written in C++ for simulation and reconstruction. To simulate the detector response, AliRoot transports to other packages like Geant3 [182], and Geant4 [183].

The reconstructed data, both experimental and simulated, are stored in files in an ESD (Event Summary Data) format, containing all the physical information needed for the analysis and checking the reconstruction quality for each sub-detector. There is another data format, an AOD (Analysis Object Data) format, which contains only the information needed for the analysis.

The Worldwide LHC Computing Grid (WLCG) is used for simulation creation, data reconstruction, and analysis [184]. The Grid is a network of computing facilities from more than 40 different nations that is geographically distributed. These data centers are organized into three tiers, demonstrating a hierarchy. Tier 0 computing facilities, located at CERN and Budapest's Wigner Research Centre for Physics, store copied data and execute raw data reconstruction. The Tier 1 centers, which are

also used to reprocess the data and carry out a portion of the reconstruction, hold a second copy. Tier 2 centers concentrate on the analysis and production of Monte Carlo simulations.

AliRoot also allows to access the grid via AliEn (ALICE environment). [185]. The AliEn user interface helps to access data, send analysis tasks to the grid, simulation generation, and job status monitoring.

3.4 Vertex and track reconstruction

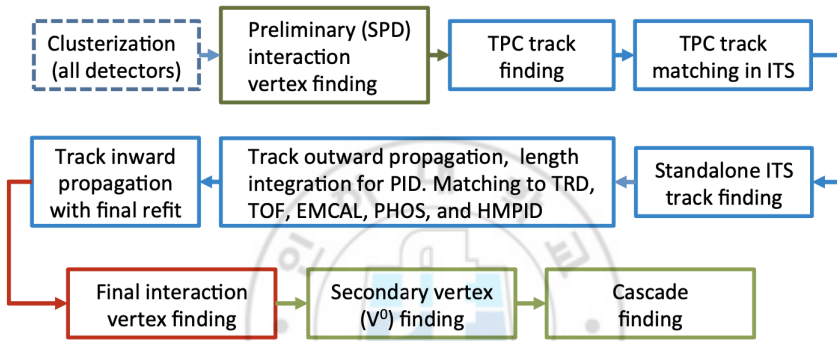


Figure 3.7: Event reconstruction flow. Figure from [31].

After data acquisition, ALICE performed an event reconstruction, as depicted in Fig. 3.7. The raw data are related to signals generated within the detector by the same interaction. The subsequent reconstruction is identical to the Monte Carlo simulation and real data. The steps for performing vertex and track reconstruction are as follows:

Clusterization The detector data are converted into clusters characterized by positions, signal sizes, signal times, and associated errors. This process is performed individually for each detector.

Track Reconstruction Using clusters from the first two ITS layers, SPD, the major interaction vertex is estimated. With the Kalman filter technique [186], track reconstruction and fitting are carried out using the TPC and ITS information. The other center barrel detectors are matched to and fitted with the tracks that the ITS-TPC detectors have rebuilt. Using the reconstructed tracks, the final interaction vertex is established.

Secondary Vertex Finding If the primary vertex and all tracks are reconstructed, ALICE can reconstruct the secondary vertex generated by strange hadrons.

3.4.1 Preliminary determination of the primary vertex

A preliminary decision of the interaction vertex, connecting two clusters on the SPD, one cluster per layer, constitutes the first SPD track reconstruction stage. The space points where numerous tracklets converge are explored for candidate vertices. The interaction vertex is determined to be the vertex, which is the place at which all maximal tracklets converge.

3.4.2 Track reconstruction

The track finding and fitting procedure are described following the inward-outward-inward scheme shown in Fig. 3.8.

TPC Track Reconstruction Finding tracks in the TPC at a large radius is the first step of the tracking algorithm. Track seeds are constructed first with two TPC clusters and the vertex point (SPD Vertex), and then with three clusters without the vertex constraint. At each step, the proximity cut is used and track seeds are propagated inward. Only tracks that miss fewer than 50% of the expected clusters for a track point and have at least 20 clusters (out of a maximum of 159 available) are considered.

ITS-TPC Track Reconstruction The hits in the outermost SSD layer match to the rebuilt TPC tracks. The new seeds are then updated at each ITS layer by all clusters within a proximity cut, after which the seeds are propagated inward and updated. In the absence of the extrapolation occurring in the dead zone of the layer, where there is no cluster, a χ^2 is taken into account whenever an extrapolated seed fails to locate a compatible cluster. Each TPC track thus generates a tree of track hypotheses in the ITS. The contenders are chosen for the ITS+TPC track in the rebuilt event based on their χ^2 scores.

Re-fitting and Finalisation The ITS-TPC tracks are back-propagated outwards to match the clusters in the TRD, TOF, and the signals in other detectors outside of TOF, as shown in the 2nd path of Fig. 3.8. Then, all the tracks are propagated back to the innermost ITS layer with the Kalman Filter refit pass, using the information from the attached clusters, and completing the three passes to fit the track parameters. Fig. 3.9 shows the resolution on $1/p_T$ for tracks of TPC tracks and ITS-TPC tracks

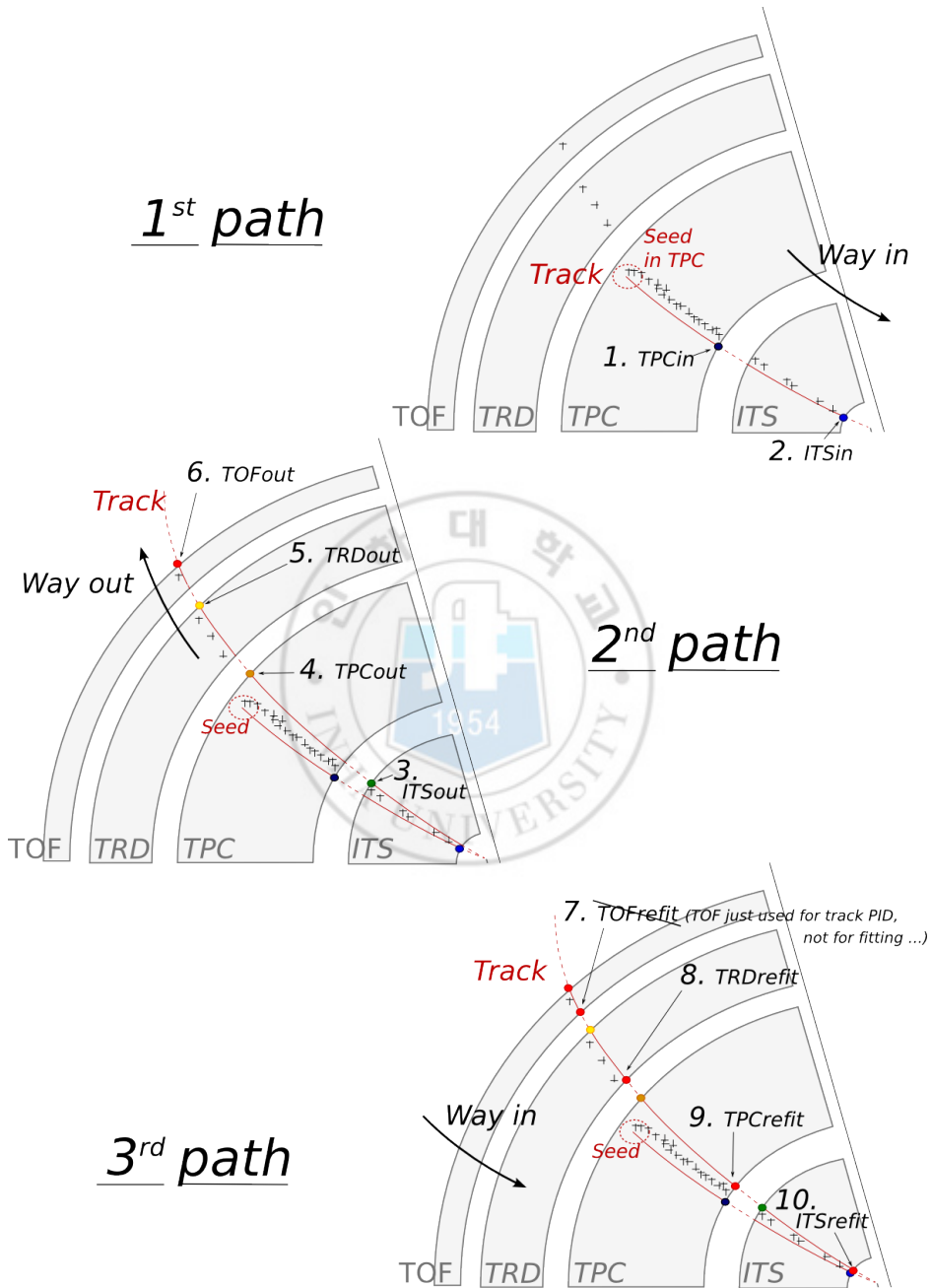


Figure 3.8: Tracking algorithm in ALICE. Figure from [33].

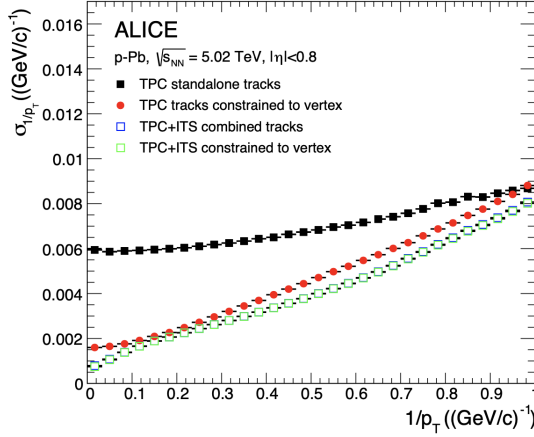


Figure 3.9: The p_T resolution for TPC and ITS-TPC combined tracks with and without constraint to the vertex. The constraint of the vertex improves the resolution of TPC standalone tracks. However, for ITS-TPC tracks, it has no effect (green and blue squares overlap). Figure from [34].

with and without vertex constrain in p-Pb $\sqrt{s_{NN}} = 5.02$ TeV. The $1/p_T$ resolution of the track related to the p_T resolution is:

$$\frac{\sigma_{p_T}}{p_T} = \frac{\sigma_{1/p_T}}{1/p_T} \quad (3.3)$$

The covariance matrix of the fit of the Kalman filter produces the $1/p_T$. The resolutions increase from 1% to 10% in the p_T range between 0.1 and 100 GeV/c.

3.4.3 Final determination of the interaction vertex

A re-determination of the interaction vertex takes over the one performed with the SPD tracklets when the final set of global tracks is decided, described in 3.4.1. The global tracks are propagated inward to their point of closest approach, and a first determination of the interaction vertex takes place. The far outliers with respect to the nominal beam pipe are ignored. Then, a more precise vertex fitting is proceeded using the track-weighting procedure described in [187] to reject the remaining outliers' contributions.

In Fig. 3.10, the resolution of the final interaction vertex position as a function of the charged particle multiplicity in pp collisions at $\sqrt{s} = 7$ TeV is shown. The precision of the determination for vertex position with the global tracks rapidly increases

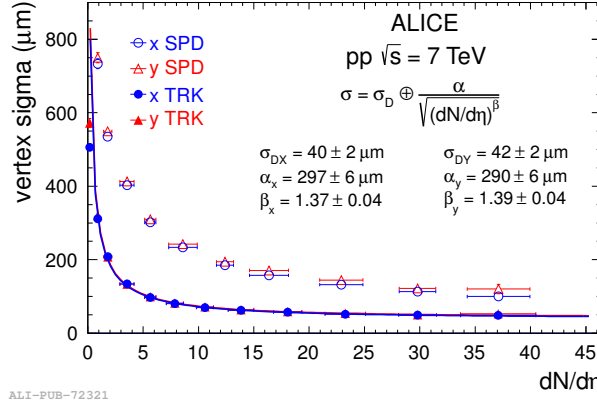


Figure 3.10: Transverse width of the final interaction vertex distribution (solid points), consist of the luminous region σ and the vertex resolution $\alpha/\sqrt{(dN/d\eta)^\beta}$. The widths of the preliminary (SPD) interaction vertices are shown (open points). Figure from [34].

as a function of multiplicity, and its trend can be parametrized with the power law function, $\alpha/\sqrt{(dN/d\eta)^\beta}$. The primary vertex determination improves more than a factor 2 in resolution with respect to the algorithm mentioned in Sec. 3.4.1.

3.4.4 Secondary vertex finding using weak decay topology

ALICE reconstructs secondary vertices from particles that either have a long decay length or interact with the detector's material. The decay topology of K_s^0 , and Ξ^- is shown in Fig. 3.11. The first tracks chosen are those whose Distance of Closest Approach (DCA) to the interaction vertex is greater than a predetermined minimum value (0.5 mm in pp and 1 mm in Pb—Pb). The Point of Closest Approach (PCA) between the two selected tracks, called V0 candidates, is determined for each pair of unlike-sign tracks. The V0 candidates are identified to further the following cuts:

- Less than 1.5 cm is the distance between the two tracks from their PCA.
- The inner-most hit of either of the two tracks is far from the interaction vertex than the PCA.
- The total momentum vector of the pair and the line connecting the primary (interaction) and secondary vertices have a Cosine of the Ponting Angle (CPA) θ that is greater than 0.9.

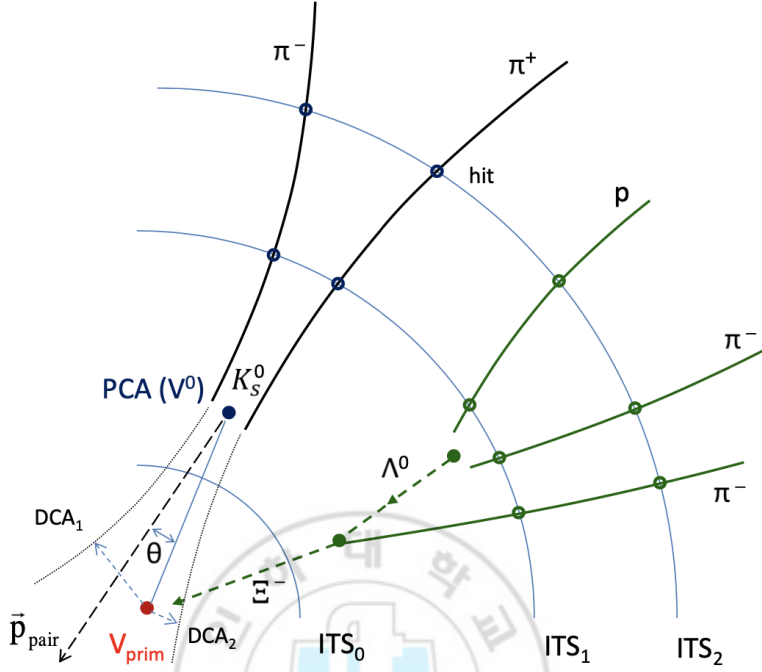


Figure 3.11: Secondary vertex reconstruction algorithm, showing the K_s^0 and Ξ^- decays. The decay points were placed between the first two ITS layers. The solid lines indicate the reconstructed charged particle tracks extrapolated to the secondary vertex. Extrapolations to the primary vertex and auxiliary vectors are shown with dashed lines. Figure from [34].

The invariant mass of K_s^0 (left) and Λ (center) in central Pb—Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV are shown in Fig. 3.12. The energy loss dE/dx of pions and protons, which are daughters of Λ , was used to identify them in the TPC detector with $p_T < 1.5$ GeV/ c (see Sec. 3.5.1). The reconstruction efficiency of K_s^0 and Λ in central (0-5%) and peripheral (80-90%) collisions are presented as a function of p_T in the right panel of Fig. 3.12. Due to the lower probability of decay in the fiducial volume ($r < 100$ cm) of the V0 search at high p_T , the reconstruction efficiency of Λ is on the decrease at high p_T .

Following V0 reconstruction, the cascade (Ξ^-) reconstruction is carried out as schematically indicated in Fig. 3.11. Reducing the candidates' reciprocal distance at the PCA, a secondary track (π^-) is matched with a V0 candidate. The latter must be outside of a cylinder volume centered on the interaction vertex ($r > 0.2$ cm) in

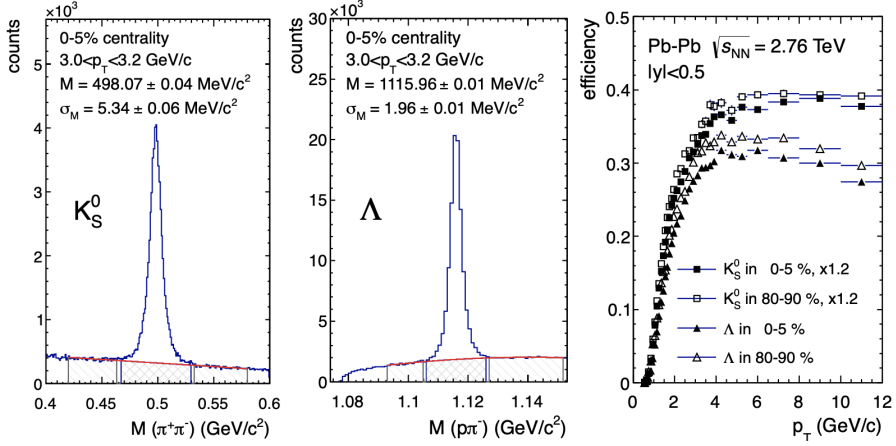


Figure 3.12: Left: invariant mass distributions of $\pi^+\pi^-$ pairs Middle: invariant mass distributions of $p\pi^-$ pairs The hatched areas show the regions of the K_S^0 and Λ peaks and the combinatorial background in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. Right: reconstruction efficiencies as a function of p_T for central (0-5%) and peripheral (80-90%) collisions. Figure from [34].

order for these candidates to be considered.

These V0 and cascade candidates are included in the data for analysis and are used analysis described in Ch. 4.

3.5 Particle identification

3.5.1 Particle identification in the TPC

The TPC is important for identifying charged particles over a wide momentum range. The PID is performed by measuring the ionization charge signal produced by a charged particle on the 159 pad rows of the TPC end plates, as mentioned in Sec. 3.2.2. The final estimate of the produced energy by the single track consists of a truncated average of 65% of lowest-amplitude signals to reject the secondary ionization contributions. By measuring this amount for each track and reconstructing the momentum p , particle identification is performed by comparing the measured dE/dx value with the expected value for each particle with momentum p . The expected dE/dx for each particle is obtained with a data-driven approach, using the

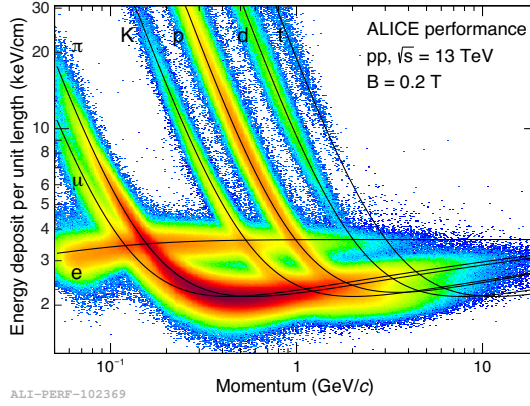


Figure 3.13: Charged particles dE/dx as a function of momentum measured by TPC in pp collisions at $\sqrt{s} = 13$ TeV. Figures from ALICE Figure Repository.

following parametrization based on the Bethe-Block equation:

$$f(\beta\gamma) = \frac{P_1}{\beta^{P_4}} [P_2 - \beta^{P_4} - \ln(P_3 + \frac{1}{(\beta\gamma)^{P_5}})] \quad (3.4)$$

where $P_1 - P_5$ are free parameters, β is the relative velocity of the particle to the light, and $\gamma = 1/\sqrt{1 - \beta^2}$ is the Lorentz factor.

The measured dE/dx in the TPC as a function of momentum in pp collisions at $\sqrt{s} = 13$ TeV is shown in the Fig. 3.13. The solid lines are the expected trends for different particle species, according to the Eq. 3.4. The largest separation is shown at low momentum region ($p < 0.7$ GeV/c), where the particle can be identified on a track-by-track basis. At higher momentum region, where a nearly constant separation among the several species is visible on a wide momentum range, particle identification can still be possible with a statistical approach based on the multi-Gaussian fit of the dE/dx distribution. This performance is shown an overall good dE/dx resolution, which amounts to 5.2% in pp collisions [34].

3.5.2 Particle identification in the TOF

TOF identifies charged particles by measuring the time of flight on a fixed length L . The time of flight can be obtained with a particle with mass m and knowing its momentum p :

$$t = \frac{L}{v} = L \cdot \frac{E}{pc^2} = L \cdot \frac{\sqrt{m^2 p^2 + m^2 p^4}}{pc^2} \quad (3.5)$$

where E is the energy of the particle, and c is the speed of light.

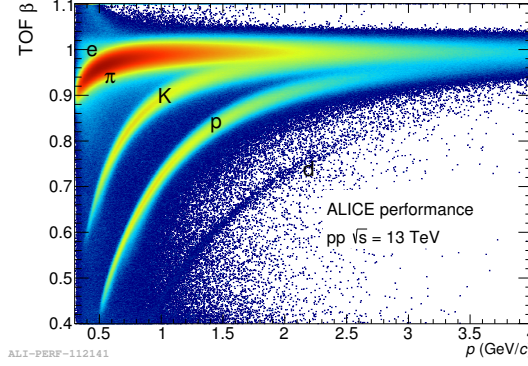


Figure 3.14: β of charged particles as a function of momentum measured by TOF in pp collisions at $\sqrt{s} = 13$ TeV. Figure from ALICE Figure Repository.

The velocity of charged particles measured by the TOF in pp collisions at $\sqrt{s} = 13$ TeV is shown in Fig. 3.14. TOF helps the particle identification in the intermediate momentum region, ensuring a kaon versus pion separation better than 3σ up to $p \simeq 2.5$ GeV/ c , and a kaon versus proton separation until $p \simeq 4.5$ GeV/ c .

Chapter 4

Ξ_c^0 production in pp collisions at $\sqrt{s} = 13 \text{ TeV}$

The ALICE detector at the Large Hadron Collider(LHC) is optimized to investigate the Quark-Gluon Plasma (QGP) produced in heavy-ion collisions. Charm quark is a sensitive probe to clarify the properties of the QGP, thanks to its large mass. They are mainly produced at the initial stage of the collisions via hard scattering processes and experience the whole evolution of the system. Measurements of charm-baryon production in pp collisions are essential to be a reference for Pb–Pb collisions studies. Furthermore, the measurements in pp collisions provide the constraint of pQCD calculations and for model calculations of charm hadronization in a vacuum. The charmed baryon-to-meson ratio can provide unique insight into hadronization processes.

In this chapter, prompt Ξ_c^0 baryon production decaying into $\Xi_c^0 \rightarrow e\Xi\nu$ in pp collisions at $\sqrt{s} = 13 \text{ TeV}$ is presented, which is one of the charmed baryons.

$$\Xi_c^0 \rightarrow e^+ \Xi^- \nu_e \text{ and charge conj.}$$

The data, Monte Carlo sample, and trigger condition are introduced in Sec. 4.1. The procedure of raw yield extraction is described in Sec. 4.2, and Monte Carlo correction for converting the raw yield to cross section is described in Sec. 4.3. Details of the systematic uncertainties study are described in Sec. 4.4. In Sec. 4.5, The production of prompt Ξ_c^0 and charmed baryon-to-meson and the charmed baryon-to-baryon ratio are reported. Additionally, the branching fraction ratio for two decay channels of Ξ_c^0 , $\text{BR}(\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e) / \text{BR}(\Xi_c^0 \rightarrow \Xi^- \pi^+)$ ratio is reported.

4.1 Data sample and event selection

4.1.1 Data sample

The minimum bias (MB) triggered data sample from pp collisions collected at the LHC using ALICE detectors in 2016, 2017, and 2018 with a center-of-mass energy of $\sqrt{s} = 13$ TeV used. The detail of MB trigger is introduced in Sec. 4.1.2.

The Monte Carlo (MC) sample was generated using PYTHIA8, which is $c\bar{c}$ pair was generated for each event, and Ξ_c^0 was used as trigger particles to store the events. In MC sample, the Ξ_c^0 was forced to enhance the statistics and to decay in the semi-leptonic decay channel ($\Xi_c^0 \rightarrow e^+ \Xi^- \nu_e$ and charge conj.) for each event. The number of events in the MC sample was 50 millions.

4.1.2 Event selection

Events were selected with the Minimum Bias (MB) trigger, which is the trigger configuration that imposes the smallest bias by accepting most of the events. Pile-up events, which have more than one primary vertex in the reconstructed event, were not included. The position of the primary vertex was correlated with the efficiency of the z-Vertex reconstruction, The efficiency would be approximately independent of the vertex position within the selected range of $|v_z| < 10$ cm. The event selection was implemented to the data and MC samples. The number of events of data after the event selection was 1862M, and the number of events of MC after the event selection was 42M for pp collisions at $\sqrt{s} = 13$ TeV, considering the vertex efficiency.

The summary of event selection is listed in the following contents:

- Trigger Class: MB trigger
- Rejection of pileup events only for the real data
- Primary vertex reconstructed from ITS+TPC tracks within ± 10 cm from the nominal interaction

4.2 Raw yield extraction

The production of Ξ_c^0 decaying to $\Xi_c^0 \rightarrow e\Xi\nu$ is introduced in this section. The electron and Ξ were paired to reconstruct the Ξ_c^0 candidates. The identification of electrons is described in Sec. 4.2.1, the reconstruction of Ξ is described and in Sec. 4.2.2. The neutrino cannot be detected because there is no interaction with detectors. Therefore, the kinematics of neutrino is missing. Therefore, the unfolding technique was used to correct the missing neutrino kinematics in $e\Xi$ pair, as described in Sec. 4.2.6.

4.2.1 Electron identification

Electrons selected by the ALICE detectors with the capabilities of particle identification (PID) were taken into account to reconstruct the Ξ_c^0 candidates.

The pseudorapidity, which is defined as $\eta = -\ln(\tan(\theta/2))$ with θ the angle between a particle's track and the beam axis, was limited to $-0.8 < \eta < 0.8$. To get the pure electrons, the maximal Distance of the Closest Approach (DCA) of the tracks to the collision vertex was restricted to 2 cm for the beam direction (z -direction) and 1cm for the radial direction (xy -direction), which are the loose DCA cuts.

The ALICE detectors have several different subsystems to identify hadrons and electrons. Because each detector's limited acceptance and efficiency in ALICE are limited, it is necessary to set the usage of different detectors. Therefore, two different PID detector systems, Time-of-Flight (TOF) and Time-Projection-Chamber (TPC) (see Sec. 3.2.3 and Sec. 3.2.2), were used together to get a higher purity of the electrons.

The TOF was used to reduce the amount of hadron contamination. Tracks inside three times the standard deviation (σ) of the electron's expected time-of-flight were used. This led to a suppression of pions, kaons, and protons at low momenta. The energy loss is defined as the parametrized electron Bethe-Bloch line deviation divided by the energy loss resolution (called TPC $n\sigma$). After $3\sigma_{\text{TOF}}$ cuts, the distribution of $n\sigma_{\text{TPC}}$ as a function of p_T is shown in the Fig. 4.1.

In Fig. 4.1, the straight line represents the upper limit of electron selection, and the curved line represents the lower limit of electron selection. The lower limit of electron selection was defined as following the previous analysis strategy [38]. Electrons were selected within the upper and lower lines in the first panel of Fig. 4.1. The others in Fig. 4.1 show $n\sigma_{\text{TPC}}$ distributions in different p_T intervals, projected from the first panel of Fig. 4.1. The $n\sigma_{\text{TPC}}$ distributions were fitted with the Gaussian function for pion (π), colored as blue, and for the electron (e), colored as red, and

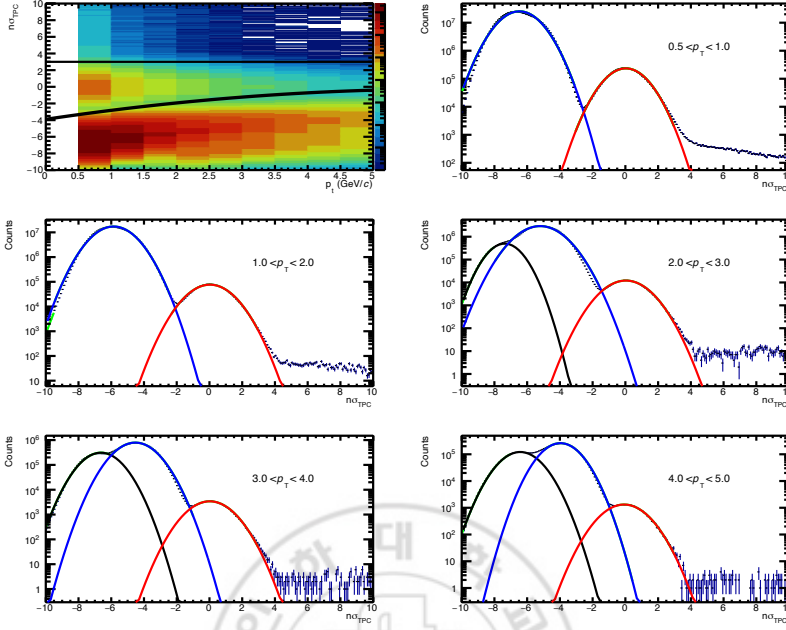


Figure 4.1: First panel: the distribution of $n\sigma_{\text{TPC}}$ as a function of p_T . The straight line represents the upper limit of electron selection, and the curved line represents the lower limit. Others: the distribution of $n\sigma_{\text{TPC}}$ for different electron's p_T bin projected from the first panel. The distribution of the red (blue, black) line is the contribution of electrons (pions, proton, or kaon).

for kaon (K) and proton (p), colored as black. As p_T increases, it is not easy to select electrons due to the overlap of the contribution from each particle.

The p_T -dependent selection of lower limit for $n\sigma_{\text{TPC}}$ was determined by comparing the purity and contamination with other selections. The electron purity and contamination were estimated by integrating the $n\sigma_{\text{TPC}}$ distribution, as below:

$$\text{Purity} = \frac{\int_{x_{\min}}^{x_{\max}} f_{\text{elec}}(x) dx}{\int_{x_{\min}}^{x_{\max}} (f_{\pi}(x) + f_{p/k}(x) + f_{\text{elec}}(x)) dx} \quad (4.1)$$

$$\text{Contamination} = \frac{\int_{x_{\min}}^{x_{\max}} (f_{\pi}(x) + f_{p/k}(x)) dx}{\int_{x_{\min}}^{x_{\max}} (f_{\pi}(x) + f_{p/k}(x) + f_{\text{elec}}(x)) dx} \quad (4.2)$$

where the $x_{\min}$ and $x_{\max}$ are the lower and upper limit of the $n\sigma_{\text{TPC}}$ cut value used for the electron identification.

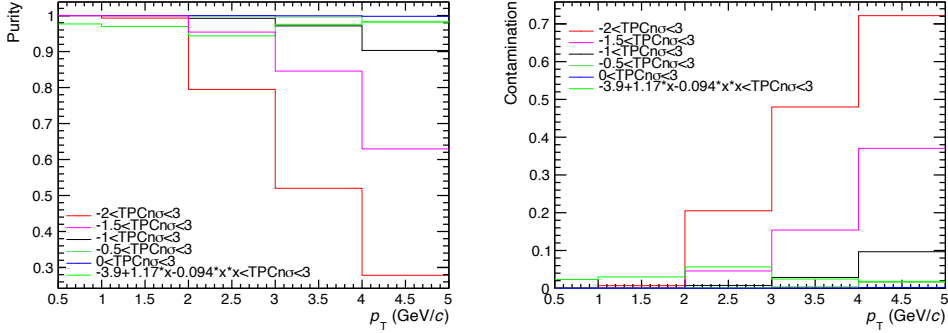


Figure 4.2: Left: The purity of electrons for varied selections. Right: the contamination of electrons for varied selections.

Cuts variables	cuts
AOD Filter Bit	4
$ n\sigma_{TPC} $	< 5

Table 4.1: Loose election selections for determination of prefilter selection.

The purity and contamination of electrons with different $n\sigma_{TPC}$ selections as a function of p_T are shown in Fig. 4.2. The green line, which means the p_T -dependent selection showed a better description for electron purity and contamination, comparing other variations.

The dominant background of electrons was mostly from Dalitz decays ($\pi^0 \rightarrow \gamma e^+ e^-$) of light neutral mesons like π^0 or η mesons or photon conversions in the material. They have the opposite charge, which is *photonic decays*:

$$\begin{aligned}\eta &\rightarrow e^+ e^- \gamma \\ \pi^0 &\rightarrow e^+ e^- \gamma \\ \gamma &\rightarrow e^+ e^-\end{aligned}$$

To suppress electrons from photon conversions in the ITS hits in both of the innermost SPD layers (the two innermost layers of ITS) were required for all selected tracks. Furthermore, one additional hit was required in the other four layers of the ITS to reject electrons paired with the same charge, which usually does not originate from the same decay.

The electron candidates are paired with opposite-sign electron candidates passing very loose electron selection (Table 4.1) in the same event. The pairs with the

minimum invariant mass were selected to maximize the probability of finding the electron partner from a photonic decay.

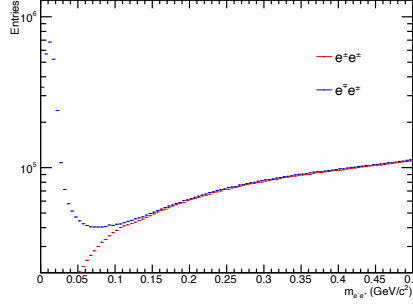


Figure 4.3: Minimum invariant mass distribution of electron pairs.

The invariant mass distribution of pairs shown in Fig. 4.3 has a peak at $m_{ee} \sim 0 \text{ GeV}/c^2$, which clearly shows the contributions from the photonic decays. The $m_{e^+e^-} > 0.05 \text{ GeV}/c^2$ selection was applied to reject the photonic electrons, called as *Prefilter*. However, the signal electrons having a partner that in the combinatorics $m_{e^+e^-} < 0.05 \text{ GeV}/c^2$ can be rejected accidentally. Such efficiency loss can be estimated using the same-sign electron pairs because they are totally combinatorial pairs. The efficiency loss of the selection turned out to be very small. A detail of prefilter efficiency is in Sec. 4.2.4.

The summary of electron selections is shown in Tab. 4.2.

Cuts vairbales	Cuts
AOD Filter Bit	4 (Standard cuts with very loose DCA)
Number of TPC crossed rows	> 70
Number of TPC crossed rows/Number of findable cluster	> 0.8
Number of TPC PID clusters	> 50
ITS/TPC refit	TRUE
Number of ITS cluster	≥ 3
$p_T \text{ (GeV}/c)$	> 0.5
$ \eta $	< 0.8
SPD hit	both
$ n\sigma_{TOF} $	< 3
$ n\sigma_{TPC} $	$> -3.9 + 1.2 \cdot x - 0.094 \cdot x^2, < 3$
prefilter cut	$m_{e^+e^-} > 0.05 \text{ GeV}/c^2$

Table 4.2: Selections for electron identification.

4.2.2 Ξ reconstruction

The Ξ^- baryons were reconstructed from the decay chain $\Xi^- \rightarrow \pi^- \Lambda$, followed by $\Lambda \rightarrow p \pi^-$. The characteristics of Ξ^- are shown in Table 4.3 and the decay topologies are shown in Fig. 4.4

Particle	Quarks	S	$I(J^P)$	Mass(MeV/c ²)	Decay particles	B.R.(%)	$c\tau$
$\Xi^-(\Xi^+)$	dss(dss)	-2 (+2)	$\frac{1}{2} (\frac{1}{2}^+)$	1321.71	$\Lambda\pi^-(\bar{\Lambda}\pi^+)$	99.9	4.91

Table 4.3: Properties of Ξ ; the valence quark content, the strangeness S, the isospin, the spin, and parity $I(J^P)$, the mass and the main decay channel, together with its branching ratio (B.R.) and decay length $c\tau$.

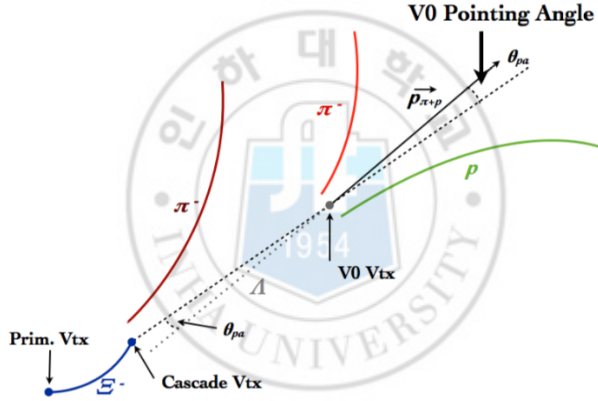
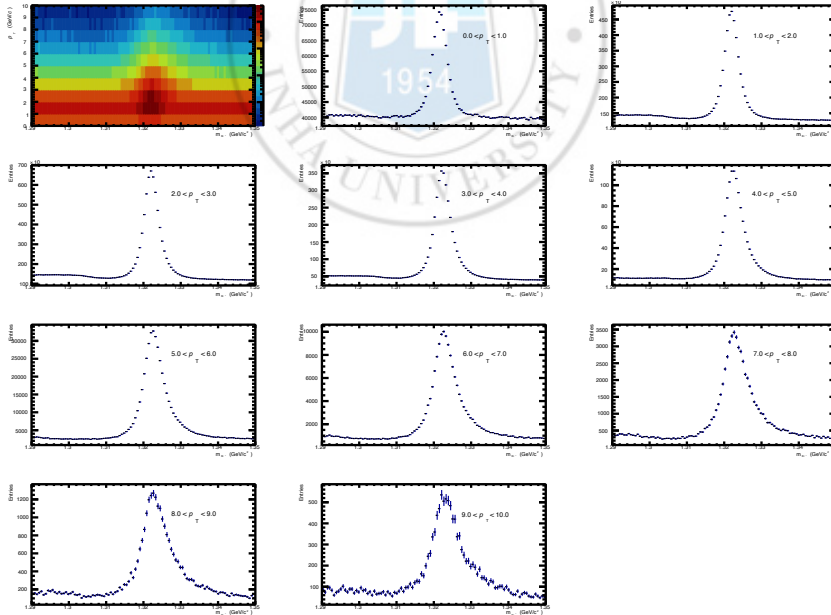


Figure 4.4: The topological variables of Ξ .

The decay daughter's track of Ξ^- used to define Ξ^- candidates were required to have the 70 of TPC crossed rows at least, Moreover, time inside four times the standard deviation of σ of the expected values for pions and protons was used. The π produced from Ξ and Λ could be identified using the different mother particles' identities: lifetime of Ξ^- and Λ , $c\tau$ is about 4.91 cm and 7.89 cm, respectively, thus they could be identified by their characteristic cascade-like and V0-shape decay topologies. The track and topology selections for Ξ are shown in Table 4.4. These selections were tuned to reduce the background while maintaining a high signal efficiency.

Fig. 4.5 shows the invariant mass distribution of Ξ^- as a function of p_T and the different p_T intervals of invariant mass distribution of Ξ^- .

Cuts variables	cuts
Number of TPC crossed rows	> 70
Number of crossed rows/Number of findable cluster	> 0.77
Λ Mass tolerance (MeV/c^2)	7.5
Ξ Mass tolerance (MeV/c^2)	8
DCA of V0 to PV (cm)	> 0.09
DCA of V0 daughters to PV (cm)	> 0.073
V0 cosine of pointing angle to Ξ vertex	> 0.983
Ξ cosine of pointing angle to PV	> 0.983
DCA of bachelor track to PV (cm)	> 0.07
V0 decay length (cm)	> 2.67
Ξ decay length (cm)	> 0.5
$ n\sigma_{TPC} $ (<i>proton</i>)	< 4
$ n\sigma_{TPC} $ (<i>pion</i>)	< 4

Table 4.4: The track and topology selections for Ξ .Figure 4.5: (Upper left) The invariant mass distribution of Ξ as a function of p_T . (The others) The invariant mass distribution of Ξ^- for different p_T intervals.

4.2.3 Raw yield estimation

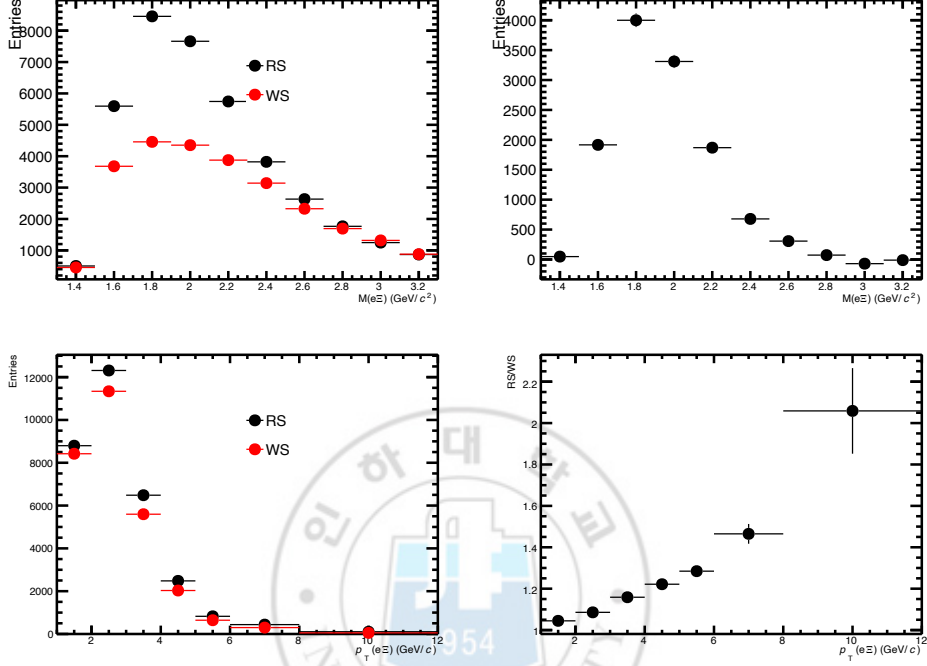


Figure 4.6: Top left: The invariant mass distributions of the Right-Sign (RS) and Wrong-Sign (WS) $e\Xi$ pairs. Top right: the WS subtracted spectra from RS, which means raw yield spectra. Bottom left: the p_T distributions of the RS and WS $e\Xi$ pairs. Bottom right: the ratio between RS and WS spectra as a function of $p_T^{e\Xi}$.

The $e\Xi$ pairs were subject to the invariant mass cut of $m_{e\Xi} < 2.5 \text{ GeV}/c^2$ and the opening angle cut of 90 degrees. The top left panel of Fig. 4.6 shows the invariant mass spectra of RS and WS $e\Xi$ pairs integrated over the whole p_T interval. It was assumed that most background sources contribute equally to WS and RS pairs. The yield of WS pairs can be treated as a combinatorial background. The invariant mass distribution of the $e\Xi$ pair doesn't have a peak at the Ξ_c^0 mass (PDG value: $2470.85^{+0.28}_{-0.40} \text{ MeV}/c^2$) [49], due to the kinematics of the neutrino. The p_T distribution of RS and WS $e\Xi$ pairs, where the invariant mass distribution from $1.3 m_{e\Xi} < 2.5 \text{ GeV}/c^2$ was integrated. The ratio between the RS and WS spectra as a function of $p_T^{e\Xi}$ was shown in Fig. 4.6 showing that the kind of low signal per background ratio after subtracting WS from RS in the $1 < p_T < 2 \text{ GeV}/c$ interval.

4.2.4 Prefilter correction

When electrons were identified as a non-photonic electrons, the prefilter selection was used (see Sec. 4.2.1). The prefilter can reject signal electrons accidentally. Therefore, a correction for the efficiency caused by the prefilter is needed. The difference between the data and MC of the prefilter is expected to depend on the event multiplicity and topology, therefore, this efficiency correction was done separately. The prefilter efficiency $\epsilon_{\text{prefilter}}$ is given by:

$$\epsilon_{\text{prefilter}} = \frac{N_{e\Xi}(\text{same-sign prefilter on})}{N_{e\Xi}(\text{prefilter off})} \quad (4.3)$$

where the $N_{e\Xi}(\text{same} - \text{sign prefilter on})$ is the number of $e\Xi$ pairs having the electron, which passed the same-sign prefilter. And $N_{e\Xi}(\text{prefilter off})$ means the number of $e\Xi$ pairs without any prefilters. The signal $e\Xi$ pair was not enough to calculate the efficiency. Therefore, the background (WS) was used to calculate the efficiency to avoid statistical fluctuations.

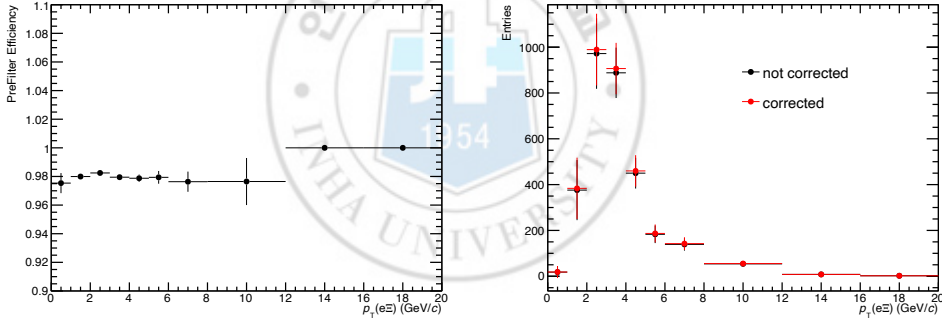


Figure 4.7: Left: $\epsilon_{\text{prefilter}}$ as a function of $p_T(e\Xi)$. Right: uncorrected spectrum of $e\Xi$ as a function of $p_T(e\Xi)$ (black), and prefilter efficiency corrected spectrum of $e\Xi$ as a function of $p_T(e\Xi)$ (red).

The prefilter efficiency as a function of $p_T(e\Xi)$ is shown in the left panel of Fig. 4.7. The efficiency loss turned out to be very small (smaller than 2%), showing the consistent loss as a function of $p_T(e\Xi)$. The prefilter efficiency corrected spectrum from the data is shown in the right panel of Fig. 4.7, colored red.

4.2.5 Over-subtraction correction

The contribution of Ξ_c^0 was extracted by subtracting the WS spectra from the RS spectra. However, in WS spectra, there were bottom baryon contributions such as $\Xi_b \rightarrow e^- \Xi^-$. The bottom baryon contributions could affect the Ξ_c^0 contributions to be underestimated. The CMS collaboration reported a measurement of the p_T -differential Λ_b^0 production cross section multiplied by the $\text{BR}(\Lambda_b^0 \rightarrow J/\psi \Lambda)$ in pp collisions at $\sqrt{s} = 7$ TeV [35], which is used for correcting the over-subtraction of the bottom baryon contributions in the WS spectra. To estimate the bottom baryon contributions in WS spectra, two assumptions were made since Ξ_b is not measured at LHC energies, and Λ_b is measured only at 7 TeV:

1. The shape of the transverse momentum distribution of $\Xi_b^{0,-}$ baryons is assumed to be the same as that of Λ_b^0 baryons.
2. The fraction of beauty quarks that hadronize into Λ_b^0 and $\Xi_b^{0,-}$ are the same as one in e^+e^- collisions.

The p_T spectrum of the Λ_b baryon was measured by CMS above 10 GeV/c [188], and the Tsallis function fitted the spectrum:

$$C \cdot p_T \left[1 + \frac{\sqrt{p_T^2 + m^2} - m}{nT} \right]^{-n} \quad (4.4)$$

where C depends on normalization, T and n are p_T shape parameters, and m is the mass of the Λ_b [35]. The shape parameters, n and T were used as $n = 7.6 \pm 0.4$ and $T = 1.1$ GeV/ c , respectively.

The Tsallis function was assumed to be valid down to $p_T = 0$, as shown in the left panel of Fig. 4.8. The ratio of the B meson cross sections at $\sqrt{s} = 13$ TeV and 7 TeV obtained with FONLL pQCD calculations, was used to scale the Λ_b^0 measurement to $\sqrt{s} = 13$ TeV. After scaling, the Ξ_b spectrum at $\sqrt{s} = 13$ TeV could be obtained by multiplying the branching fraction ratio 4.5 by scaled Λ_b^0 spectrum.

$$\frac{\text{BR}(b \rightarrow \Xi_b) \text{BR}(\Xi_b \rightarrow e \Xi \nu)}{\text{BR}(b \rightarrow \Lambda_b) \text{BR}(\Lambda_b \rightarrow J/\Psi \Lambda)} \quad (4.5)$$

The spectrum of $e\Xi$ pairs could be estimated using the refolded Ξ_b spectrum obtained by the response matrix of Ξ_b and $e\Xi$, described by Eq. 4.6.

$$N_{\Xi_b}^{\text{raw}} = \frac{d\sigma_{\Xi_b}^2}{dp_T dy} \cdot 2 \cdot \Delta p_T \cdot \Delta \cdot \varepsilon \cdot \mathcal{L}_{\text{int}} \quad (4.6)$$

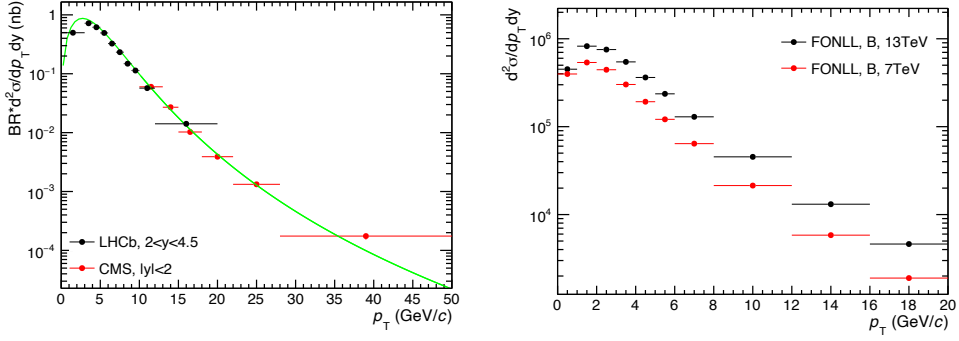


Figure 4.8: Left: Λ_b p_T spectrum at $\sqrt{s} = 7$ TeV measured by CMS [35]. Green line represents the Tsallis fit function down to $p_T = 0$. Right: B mesons p_T spectrum at $\sqrt{s} = 13$ TeV (Black) and $\sqrt{s} = 7$ TeV (Red), generated by FONLL.

where $N_{\Xi_b}^{raw}$ is the raw yield (sum of particles and antiparticles) in a given p_T interval with width Δp_T , and $\mathcal{L}_{int}$ is the integrated luminosity of pp collisions at $\sqrt{s} = 13$ TeV measured by ALICE. Factor 2 considers that the raw yield includes particles and antiparticles. The efficiency of Ξ_b , and the response matrix of Ξ_b from the MC and $e\Xi$ pairs are shown in Fig. 4.9

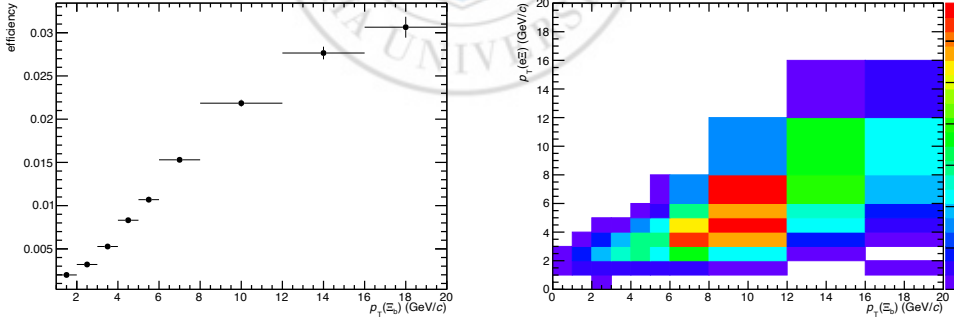


Figure 4.9: Left: Ξ_b efficiency in pp collisions at $\sqrt{s} = 13$ TeV as a function of $p_T(\Xi_b)$. Right: response matrix of Ξ_b and $e\Xi$ in pp collisions at $\sqrt{s} = 13$ TeV.

Finally, Ξ_b contribution in WS was 1 $\sim$ 2% at low p_T region, and 4% at high p_T region shown in the right panel of Fig. 4.10, due to the harder p_T spectrum of b quark compared with c quark. The $e\Xi$ pair from the bottom baryon was added to the

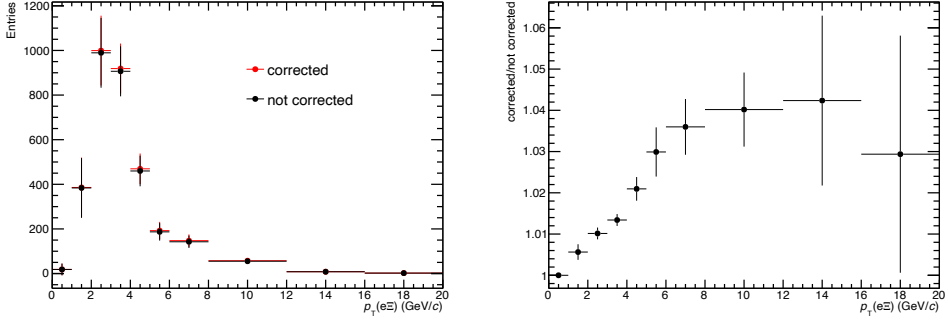


Figure 4.10: Left: the WS subtracted spectra with over-subtraction correction as a function of p_T (red), and the WS subtracted spectra without over-subtraction correction as a function of p_T (black). Right: the corrected and uncorrected $e\Xi$ spectrum ratio as a function of p_T .

$e\Xi$ pair from the raw yield spectra (RS-WS), increasing the cross section $1 \sim 6\%$.

4.2.6 Unfolding

The detector response, which can be described by a matrix R using the detector response in MC for a certain η and v_z range, is to reflect the relation between the true p_T distribution T and the measured p_T distribution M . Hence, a mathematical matrix multiplication method is needed to convert the measured yields to the true yields.

$$T = (R^{-1}) * M \quad (4.7)$$

where R , obtained from the Monte Carlo simulation, is the correlation matrix between the measured and the reconstructed p_T distribution, which also gives the conditional probability that a collision with the true distribution is measured as an event with the measured distribution.

However, R^{-1} can not be used in all cases. It would give an unreliable result with very large variances and a mathematically unbiased estimator of T . This problem can be overcome by turning it into an iterative procedure. Bayesian unfolding [189] is the default method in this analysis, which is an alternative approach to unfold the measured distribution based on the Bayes theorem, which is repeated with the unfolded distribution as a new input for the procedure until it converges.

Bayes's theorem describes the probability of an event, which is based on prior knowledge of conditions:

$$P(A|B) = \frac{P(B|A)P(A)}{P(B)} \quad (4.8)$$

$$\tilde{R}_{ji} = \frac{R_{ij}U_j}{\sum_k R_{ik}U_k} \quad (4.9)$$

where the $P(A)$ and $P(B)$ are the prior probabilities for the true distribution T and measured distribution M , respectively. And $P(B)$ is calculated with the response matrix: $M_i = \sum_k R_{ik}U_k$. U_j are the vector of the unfolded true distribution. $P(B|A)$ is the probability of measured distribution M under the condition that the true distribution T is true and is also the same with R_{ij} of the detector response matrix R . So, the $P(A|B)$ can be determined by the knowledge of the detector response matrix $P(B|A)$. Therefore, both prior probabilities of $P(A)$ and $P(B)$ are needed for the process.

$$U' = \tilde{R}[U] \cdot M \quad (4.10)$$

If spectrum U is the true distribution T , the unfolded spectrum U' is T , otherwise, it is between U and T .

The response matrix should be calculated with realistic Ξ_c^0 p_T distributions. So, the matrix was calculated in the following 3 steps:

1. The response matrix is obtained from the Monte Carlo simulation, and the Ξ_c^0 distribution is obtained with this response matrix.
2. The resulting p_T distribution of Ξ_c^0 is used to produce the response matrix for the second iteration. The p_T distribution of Ξ_c^0 from MC is corrected to the one from the first step. Then, the second p_T distribution of Ξ_c^0 is obtained with the weighted response matrix.
3. Repeat the second step until the resulting p_T distribution of Ξ_c^0 has converged.

The weighting factor was calculated as the ratio of p_T distribution of Ξ_c^0 in data and MC as shown in the left panel of Fig. 4.11. In the right panel of Fig. 4.11, the p_T distribution of measured Ξ_c^0 and the p_T distribution of Ξ_c^0 from MC were fitted with exponential functions in $p_T > 1$ GeV/ c region. The red (blue) lines obtained by maximum (minimum) uncertainty at low p_T and minimum (maximum) uncertainty at high p_T were used for systematic variation of the weighting factor.

The result of unfolding for Ξ_c^0 is shown in Fig. 4.12. The number of iterations of 3 was used in the main unfolding procedure. The measured spectrum was compared

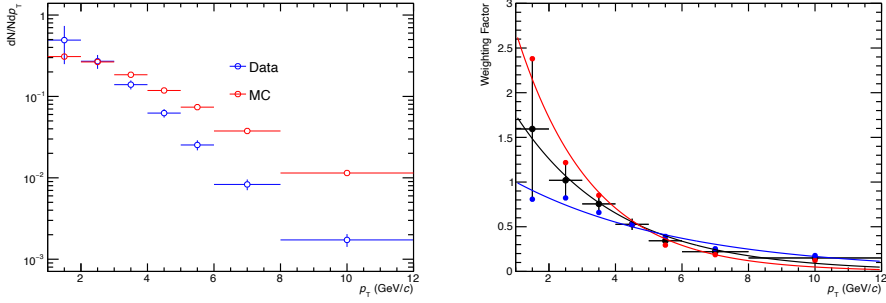


Figure 4.11: Left: the p_T distribution of Ξ_c^0 in data and Monte Carlo sample. Right: the weighting factor is defined as the ratio of p_T distribution of Ξ_c^0 in data and MC (black).

with the unfolded spectra using weighted and unweighted response matrices. Based on the refolded spectrum, the unfolding with the weighting procedure was more stable than without one.

4.3 Mote Carlo correction

4.3.1 Efficiency and acceptance correction

The total selection efficiency term, the definition of $(Acc \times \epsilon)$ as follows:

$$(Acc \times \epsilon) = \frac{N_{MC, reco}^{\Xi_c^0}}{N_{MC, gen|y|<0.5}^{\Xi_c^0}} \quad (4.11)$$

which is the total selection efficiency of $\Xi_c^0 \rightarrow e\Xi\nu$, the result is shown in Fig. 4.13.

- $N_{MC, gen}^{\Xi_c^0}$: The number of Ξ_c^0 baryons counted at the generator level in the MC simulations within Δy ($=1$), which decay into $e\Xi\nu$. The cut on the generated position of the vertex along z within 10 cm are applied to select events.
- $N_{MC, reco}^{\Xi_c^0}$: The number of Ξ_c^0 baryons counted at the reconstruction level in the MC simulations. The same analysis cuts are applied to the data. The Physics Selection and a cut on the reconstructed position of the vertex along z within 10 cm are applied to select events. Because pileup effects do not exist in the MC sample, the pileup rejection cuts are not applied in the MC analysis.

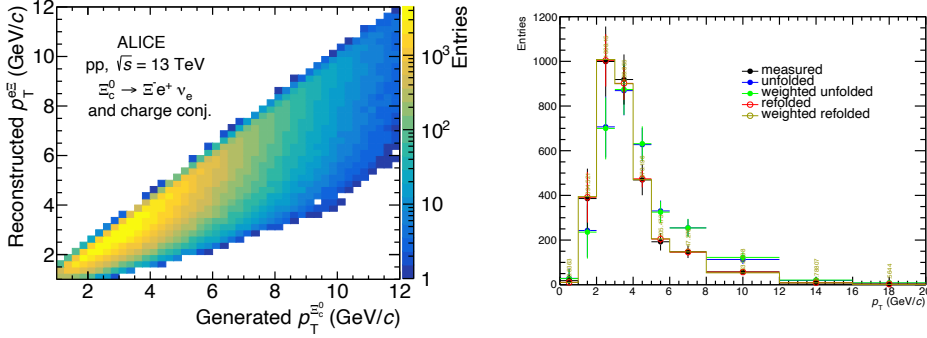


Figure 4.12: Left: the weighted response matrix of Ξ_c^0 and $e\Xi$. Right: comparison of the measured spectrum (black) with the refolded spectrum (red) and the unfolded spectrum (blue), and with the weighted refolded spectrum (yellow) and the weighted unfolded spectrum (green).

Similarly to what was mentioned in Sec. 4.2.6, the p_T distributions of the MC samples and real data are different. The exact weighting procedure was performed as in Sec. 4.2.6 to correct the difference. The $(Acc \times \epsilon)$ of the prompt Ξ_c^0 with weighting and without weighting were comparable within the uncertainties, as shown in Fig. 4.13. The $(Acc \times \epsilon)$ of prompt, feed-down, and inclusive turned out they are comparable with each other within the uncertainties. Because selections were not sensitive to the displacement by a few hundred micrometers of the prompt and feed-down Ξ_c^0 decay vertices from the collision point.

4.3.2 Prompt fraction estimation

The prompt Ξ_c^0 cross section can be obtained with the following equation:

$$\frac{d^2\sigma^{\Xi_c^0}}{dp_T dy} = \frac{1}{BR} \cdot \frac{1}{2\Delta y \Delta p_T} \cdot \frac{f_{\text{prompt}} \cdot N_{\text{raw}}^{\Xi_c^0}}{(A \times \epsilon)_{\text{prompt}}} \cdot \frac{1}{\mathcal{L}_{\text{int}}} \quad (4.12)$$

where $N_{\text{raw}}^{\Xi_c^0}$ is the raw yield, which is the sum of particles and antiparticles, in a given p_T interval in Δp_T . $(A \times \epsilon)_{\text{prompt}}$ is the product of detector acceptance and efficiency for prompt Ξ_c^0 , and $\mathcal{L}_{\text{int}}$ is the integrated luminosity, the value is $32.08 \pm 0.51 \text{ nb}^{-1}$ [190]. The Δy represents the width of the covered rapidity interval. Factor 1/2 considers that the raw yield includes particles and antiparticles, while the cross section is only for particles. The BR is the branching ratio, which is the new measurement obtained from the recently constrained absolute branching

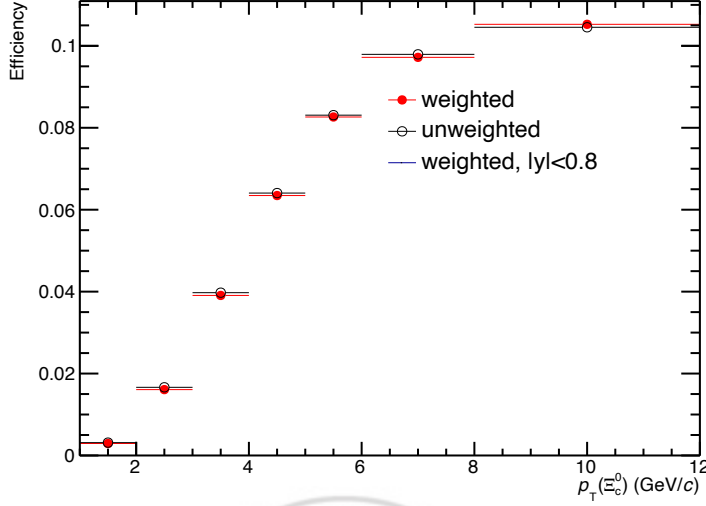


Figure 4.13: $(Acc \times \epsilon)$ of the prompt Ξ_c^0 with weighting (red) and without weighting (black) as a function of p_T .

ratio (BR) of this Ξ_c^0 decay in pp collision at $\sqrt{s} = 13$ TeV, by measuring the Ξ_c^0 production via two different decay channels, $\Xi_c^0 \rightarrow e^+ \Xi^- \nu_e$ and $\Xi_c^0 \rightarrow \pi^+ \Xi^-$ [37]. f_{prompt} is the fraction of the raw yield from prompt Ξ_c^0 , defined as:

$$f_{\text{prompt}} = 1 - \frac{N_{\text{feed-down}}^{\Xi_c^0}}{N_{\text{raw}}^{\Xi_c^0}} \quad (4.13)$$

where $N_{\text{feed-down}}^{\Xi_c^0}$ is the number of feed-down Ξ_c^0 . The $N_{\text{feed-down}}^{\Xi_c^0}$ is defined as:

$$N_{\text{feed-down}}^{\Xi_c^0} = \frac{d\sigma^{\Xi_c^0}}{dp_T dy} \cdot 2 \cdot \Delta p_T \Delta y \cdot (A \times \epsilon)_{\text{feed-down}} \cdot L_{\text{int}} \quad (4.14)$$

where $(A \times \epsilon)_{\text{feed-down}}$ is the product of detector acceptance and efficiency for feed-down Ξ_c^0 .

The production of beauty feed-down Ξ_c^0 is not measured. Therefore, a strategy based on the estimation constructed in Ref. [191] using the feed-down Λ_c^+ was adopted. The production of Λ_c^+ from Λ_b^0 -baryon decays were estimated using the b-quark cross section from FONLL calculations, multiplied by the beauty quarks fraction that fragment into Λ_b^0 . The fragmentation fraction is obtained from the beauty fragmentation function in pp collisions at $\sqrt{s} = 13$ TeV, measured by LHCb [36].

The $\Lambda_b^0 \rightarrow \Lambda_c^+ + X$ decay kinematics was modeled using PYTHIA 8.243 simulations [192], and scaled by the fraction of Ξ_b^- decaying in a final state with a Ξ_c^0 , which is taken to be 50% from the PYTHIA 8.243 generator. The contributions from other b hadrons to Ξ_c^0 were checked, which tells us that their contribution was less than 2%.

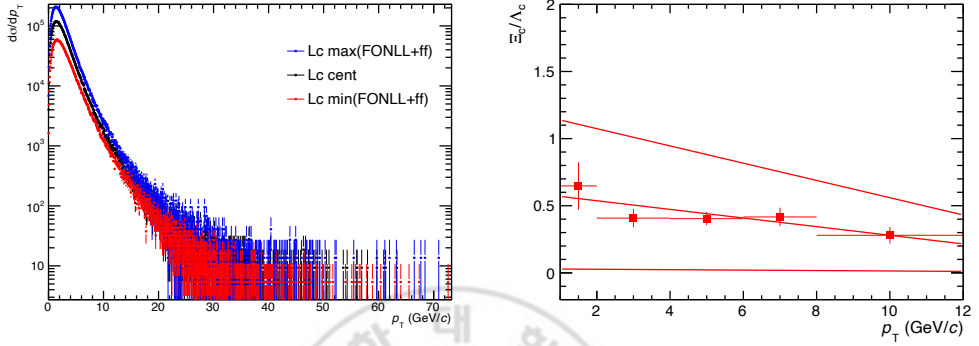


Figure 4.14: Left: The p_T distribution of feed-down Λ_c^+ obtained from the b-quark p_T -differential cross section from FONLL calculations, and the fragmentation function of beauty quark measured by LHCb [36] (Black). The p_T distribution of feed-down Λ_c^+ with the maximized variation of FONLL and fragmentation function is colored blue, and the p_T distribution of feed-down Λ_c^+ with the minimized variation of FONLL and fragmentation function is colored red. Right: the ratio inclusive Ξ_c^0 /prompt Λ_c^+ . The maximal variation and minimal variation are used to estimate the systematic uncertainties. Details are in Sec. 4.28.

The cross section of Ξ_c^0 from beauty feed-down was calculated from the cross section of Λ_c^+ originating from Λ_b^0 decays, which was scaled by the ratio of the measured p_T -differential yields of inclusive Ξ_c^0 and prompt Λ_c^+ baryons. The assumptions were needed to get a cross section of Ξ_c^0 from beauty feed-down, that the p_T shape of the cross sections of feed-down Λ_c^+ and Ξ_c^0 is the same, and that the ratio inclusive Ξ_c^0 /prompt Λ_c^+ is the same as the ratio for the feed-down baryons. The assumption was reasonable because the inclusive Λ_c^+ -baryon yield is dominated by the prompt production, based on the f_{prompt} values close to unity reported in Ref. [193].

The prompt fraction is shown in Fig. 4.15. The Ξ_c^0 yield was dominated by the prompt production, based on the f_{prompt} values close to unity.

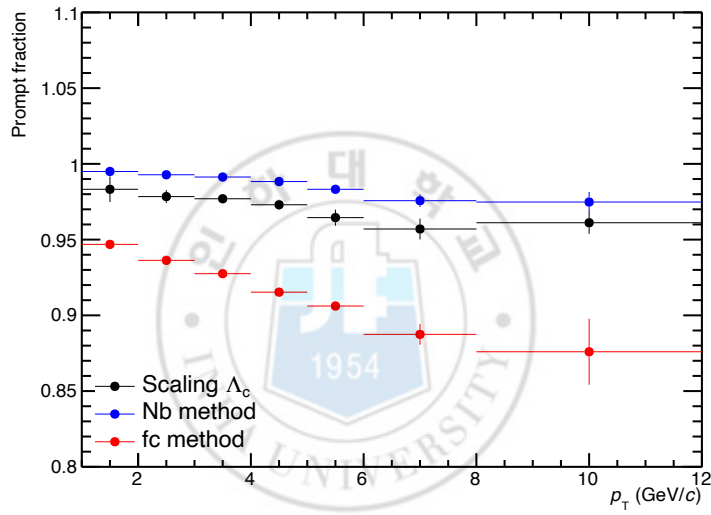


Figure 4.15: Prompt fraction of Ξ_c^0 decaying into $\Xi_c^0 \rightarrow e\Xi\nu$ in pp collisions at $\sqrt{s} = 13$ TeV.

4.4 Systematic uncertainties

The discrepancy between MC and data is caused by the imperfection of the simulation and results in systematic uncertainties. Variations of selections used to calculate systematic uncertainties show the sensitivity of the selections.

A Barlow criterion [194] has been applied for variations to prevent an effect to systematic uncertainty from statistical fluctuation.

$$B = \frac{|Default - Variation|}{\sqrt{\sigma_{Default}^2 - \sigma_{Variation}^2}} \quad (4.15)$$

If the Barlow criteria (B) is less (more) than 1, differences could result from a statistical (systematic) fluctuation.

4.4.1 Electron reconstruction

The number of TPC clusters and the number of ITS clusters were varied to study the uncertainties on the track reconstruction. Table 4.5 shows the variation for the systematic study.

Cut variables	VeryLoose	Loose	Standard	Tight	VeryTight
Number of crossed rows	> 65	> 65	> 70	> 75	> 80
Number of crossed rows/Number of findable clusters	> 0.75	> 0.8	> 0.8	> 0.85	> 0.9
Number of TPC PID cluster	> 40	> 45	> 50	> 55	> 60

Table 4.5: Electron selections used to study the uncertainties of track reconstruction. Selections that are not shown in this table are the same as the Table 4.2.

Fig. 4.16 shows the prompt Ξ_c^0 p_T spectra, the ratios of the Ξ_c^0 spectra with variations to the standard cut, efficiencies, and the RMS deviations from 1, and the Barlow values are calculated for each p_T bin. The systematic uncertainties are assigned 2%, followed by the RMS value.

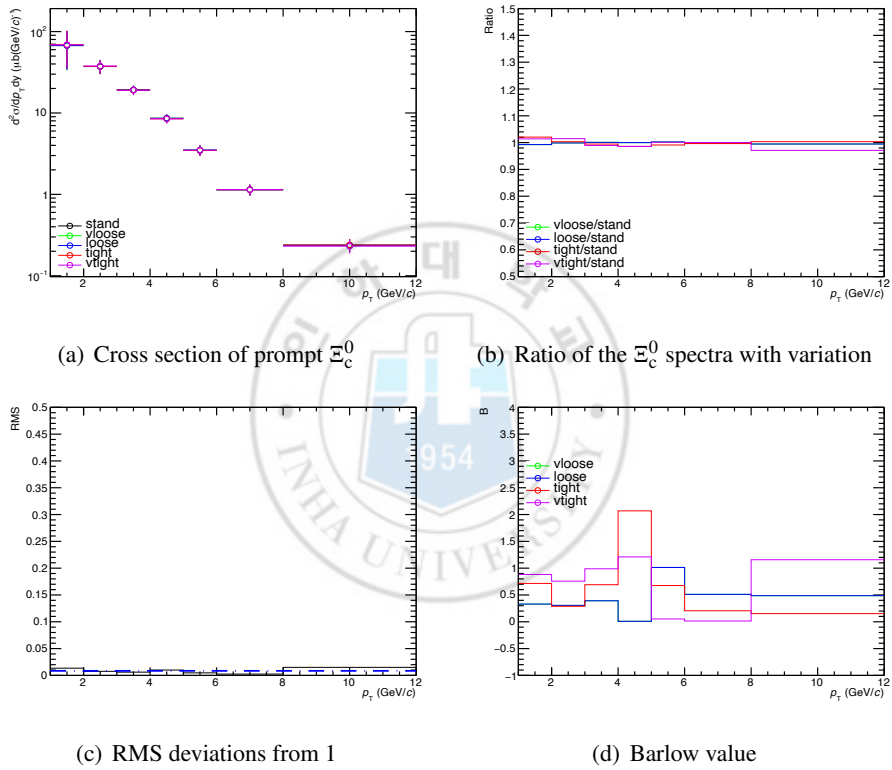


Figure 4.16: The systematic study for electron reconstruction.

4.4.2 Electron identification

The constant of the $n\sigma_{TPC}$ function criteria was varied to study the uncertainties on the track identification. The varied $n\sigma_{TPC}$ functions as a function of p_T are shown in Fig. 4.17

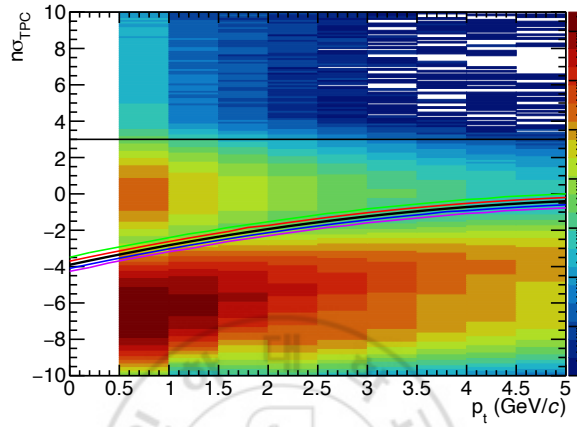


Figure 4.17: $n\sigma_{TPC}$ function variations: Standard (black), VeryLoose (violet), Loose (blue), Tight (red), VeryTight (green).

Table 4.6 shows the variation for the systematic study.

Cuts variables	VeryLoose	Loose	Standard	Tight	VeryTight
Constant of $n\sigma_{TPC}$ function	-4.3	-4.1	-3.9	-3.7	-3.5

Table 4.6: Electron selections used to study the uncertainties of track identification. Selections that are not shown in this table are the same as the Table 4.2

The systematic uncertainties are assigned as a function of p_T , followed by the RMS value.

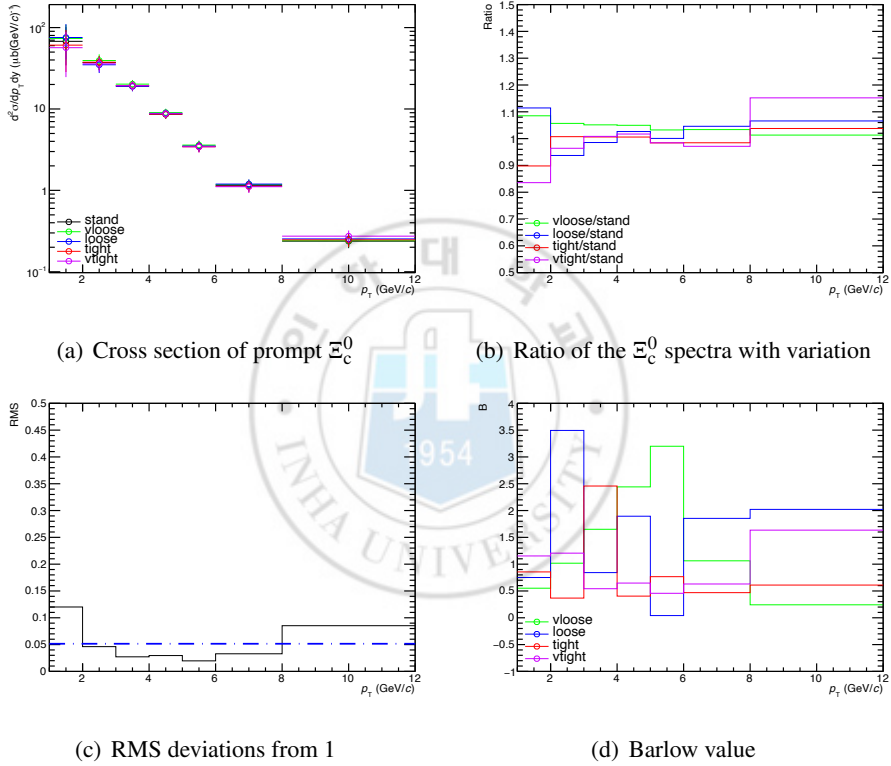


Figure 4.18: The systematic study for electron identification.

4.4.3 Ξ reconstruction

The number of TPC clusters was varied to study the uncertainties on the track reconstruction from the V0 and cascade. Table 4.7 shows the variation for the systematic study.

Cuts variables	VeryLoose	Loose	Standard	Tight	VeryTight
Number of crossed rows	> 5.96	> 66	> 70	> 70	> 74
Number of crossed rows/Number of findable clusters	> 0.74	> 0.77	> 0.77	> 0.79	> 0.79

Table 4.7: Ξ reconstruction selections used to study the uncertainties of daughter track reconstruction. Selections that are not shown in this table are the same as the Table 4.4.

The systematic uncertainties are assigned as a function of p_T , followed by the RMS value.

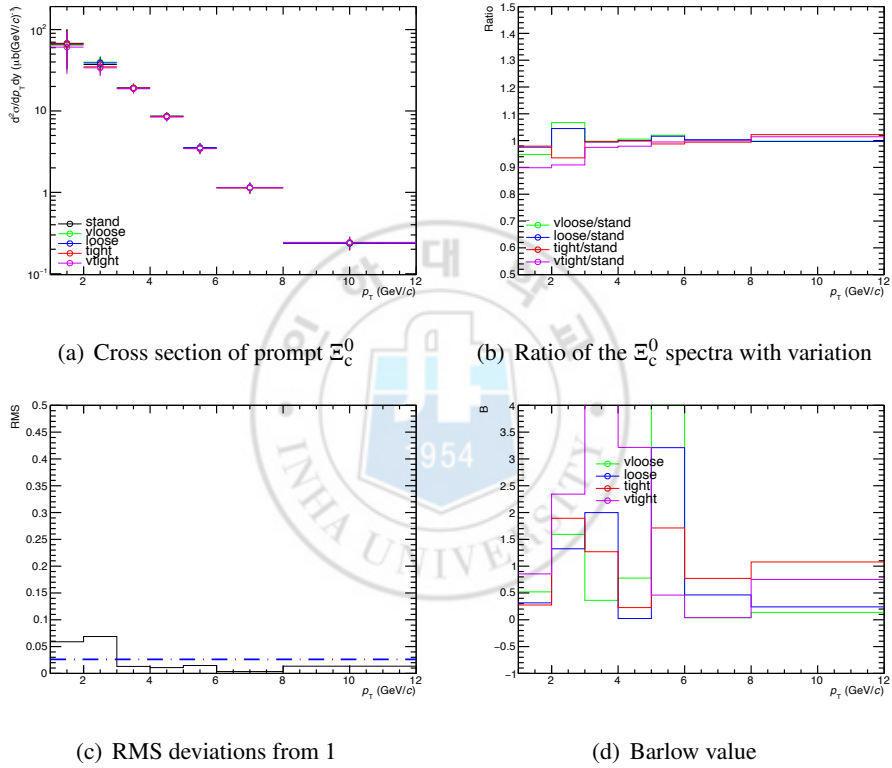
4.4.4 Ξ topology

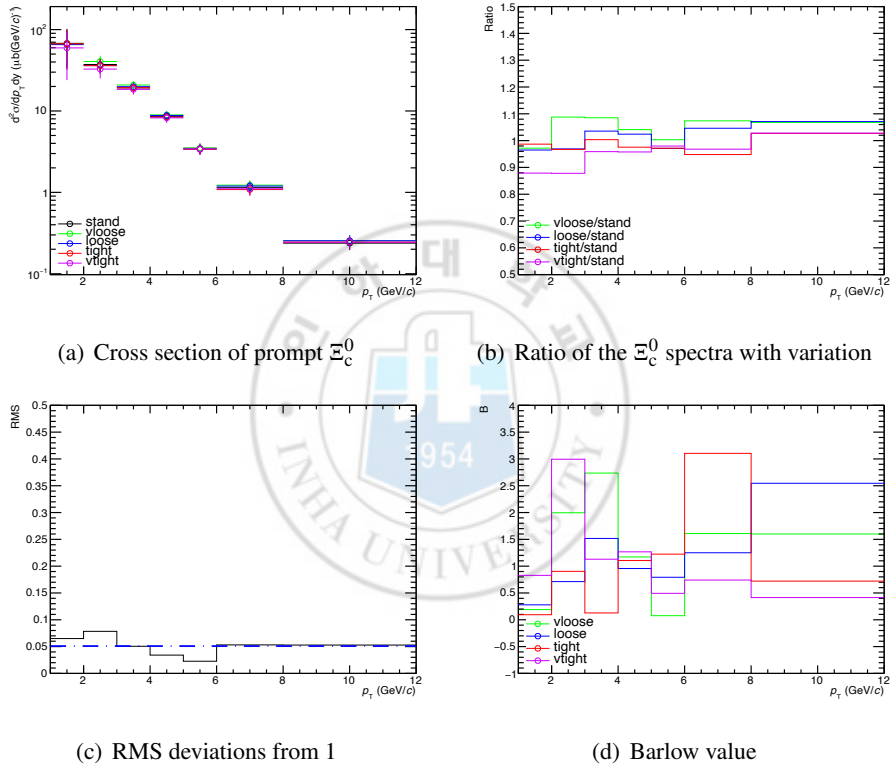
The Ξ topology selections were varied to study the uncertainties on the Ξ reconstruction from the V0 and cascade. Table 4.7 shows the variation for the systematic study.

Cut variables	VeryLoose	Loose	Standard	Tight	VertTight
Λ Mass tolerance (MeV/c^2)	7.5	7.5	7.5	7.5	7.5
Ξ Mass tolerance (MeV/c^2)	8	8	8	8	8
DCA of V0 to PV (cm)	> 0.05	> 0.07	> 0.09	> 0.11	> 0.13
DCA of V0 daughters to PV (cm)	> 0.05	> 0.061	> 0.073	> 0.088	> 0.102
V0 cosine of pointing angle to Ξ vertex	> 0.98	> 0.981	> 0.983	> 0.9839	> 0.985
DCA of bachelor track to PV (cm)	> 0.05	> 0.05	> 0.07	> 0.09	> 0.11
V0 decay length (cm)	> 1.1	> 1.55	> 2.67	> 3.6	> 4.39
Ξ decay length (cm)	> 0.5	> 0.5	> 0.5	> 0.53	> 0.72

Table 4.8: Ξ topology selections used to study the uncertainties of daughter track reconstruction. Selections that are not shown in this table are the same as the Table 4.4.

The systematic uncertainties are assigned as a function of p_T , followed by the RMS value.

Figure 4.19: The systematic study for Ξ reconstruction.

Figure 4.20: The systematic study for Ξ topology selection.

4.4.5 $e\Xi$ pair cut

The invariant mass and opening angle of $e\Xi$ were varied to study the uncertainties on the $e\Xi$ reconstruction. Due to the large background contribution, only one variation was used to estimate the systematic uncertainty of the opening angle. The Table 4.9 and Table 4.10 show the variation for the systematic study.

Cut variables	Loose	Standard	Tight
Invariant mass cut (GeV/c^2)	2.7	2.5	2.3

Table 4.9: Variations used to study the uncertainties of the invariant mass of $e\Xi$.

Cut variables	Standard	Tight
Opening angle cut (degree)	90°	70°

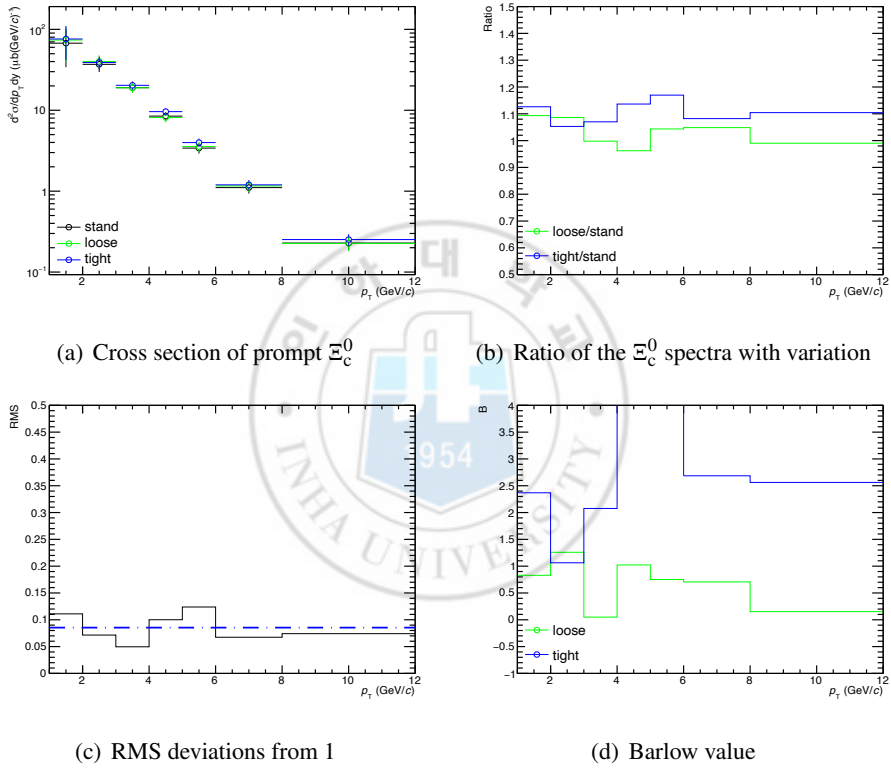
Table 4.10: Variations used to study the uncertainties of the opening angle of $e\Xi$.

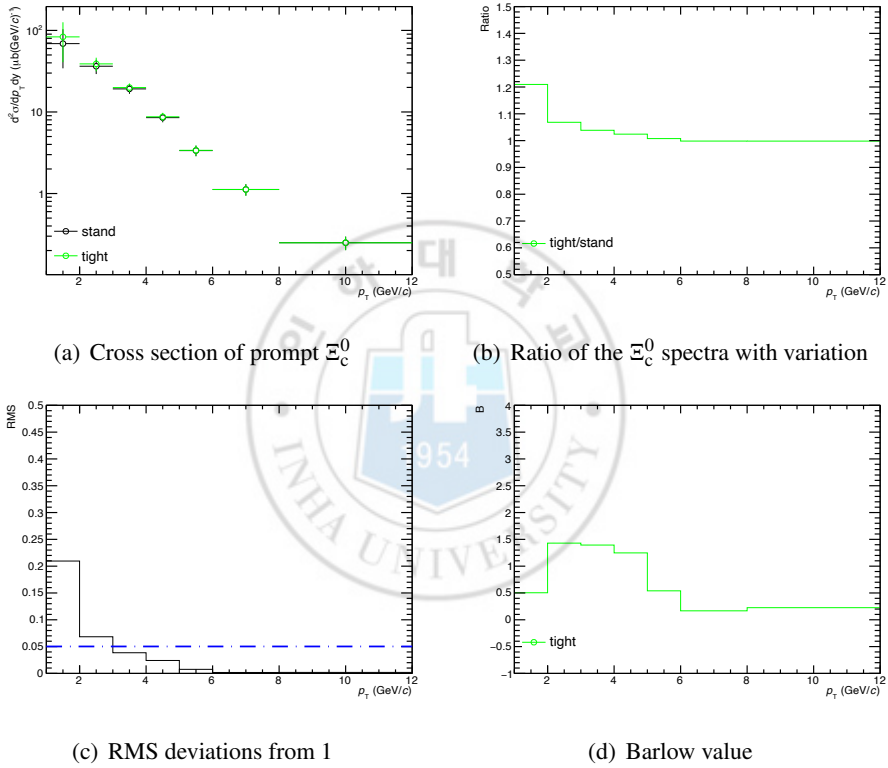
The systematic uncertainties for the invariant mass of $e\Xi$ were assigned as a function of p_T , followed by the RMS value.

The systematic uncertainties for the opening angle of $e\Xi$ were assigned as a function of p_T , followed by the RMS value. Fig. 4.22 shows that the opening angle cut is sensitive, especially at low p_T . The opening angle distributions were checked to investigate the opening angle systematic study.

The left panel of Fig. 4.23 shows the opening angle distributions of data and MC. The distributions of data for RS and WS were consistent. The red point indicates the opening angle distribution of $e\Xi$ pairs, which come from true Ξ_c^0 in whole p_T range. The green and blue points indicate that the opening angle distributions of the $e\Xi$ pairs, which come from true Ξ_c^0 in the p_T range.

If the opening angle cut becomes tighter than the standard cut (90°), the signal-to-background ratio will increase, but efficiency will be consistent with the standard cut. This effect was shown in the opening angle systematic study (see Fig. 4.22).

Figure 4.21: The systematic study for the invariant mass of $e\Xi$ pair.

Figure 4.22: The systematic study for the opening angle of $e\Xi$ pair.

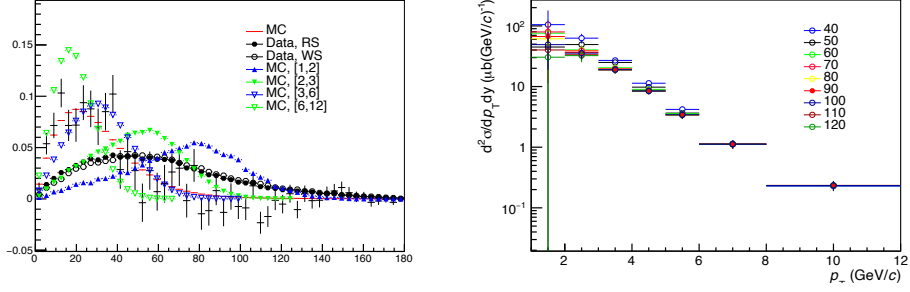


Figure 4.23: Left: the opening angle distributions of $e\Xi$ pairs. The red line indicates the distribution of $e\Xi$ pairs, which come from true Ξ_c^0 in the whole p_T range from MC, the closed black circle indicates the distribution of RS from Data, and the open black circle indicates the distribution of WS from Data. The blue and green triangle indicates the distribution from MC for each p_T interval. Right: the cross section of prompt Ξ_c^0 for each opening angle variation.

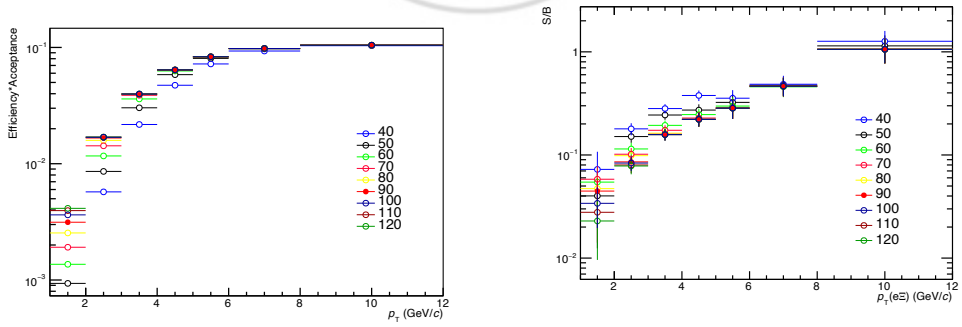


Figure 4.24: Left: The efficiencies for each opening angle variation. Right: the signal-to-background ratio for each opening angle variation.

4.4.6 Over-subtraction

The Ξ_b contributions were varied to study the systematic uncertainties on the over-subtraction.

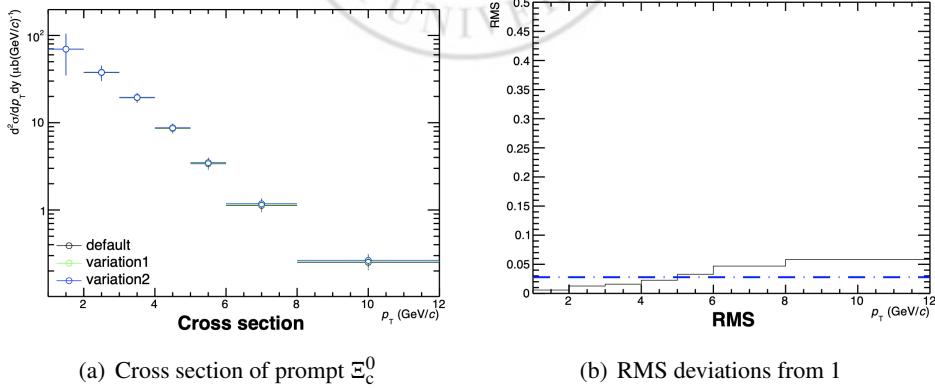
Variation 1. The uncertainties on the bottom baryons correction are estimated by varying the parameters of the TSallis functions.

- Varying n: varied by 1σ ($7.6 \rightarrow 8.0$)
- Scale up: CMS estimated the uncertainties on the cross section in the p_T range above 10 GeV/c to be $\approx 50\%$. (The normalization of the Λ_b spectrum is scaled up by the uncertainty).

Variation 2. Varying branching ratio fraction

- The uncertainties on the $B(b \rightarrow \Xi_b B)(\Xi_b \rightarrow e \Xi X)$ (30%) and $B(b \rightarrow \Lambda_b B(\Lambda_b \rightarrow J/\psi \Lambda))$ (14%) were added in quadrature. The total uncertainty of the branching fraction ratio is 33%.

Variation1 means the scale up the CMS Λ_b cross section up to 50%, and variation2 indicates the varying branching fraction ratio up to 33%. In variation1, No systematic uncertainty was assigned for varying n because systematic uncertainty was less than 1%.



(a) Cross section of prompt Ξ_c^0

(b) RMS deviations from 1

Figure 4.25: The systematic study for variations of Ξ_b contributions.

The systematic uncertainties were assigned as a function of p_T , followed by the RMS value.

4.4.7 Unfolding

p_T binning

The unfolding procedure was very sensitive to the p_T range and p_T binning, especially at low p_T . To check the p_T binning and range sensitivity of unfolding, the p_T binning and p_T range were varied. The lower limit of p_T range for the unfolded spectrum was changed from 0 (default) to 0.1 (Var1), 0.3 (Var2), and 0.5 (Var3). Also, First p_T bin was split as [0,1] (default), [0,0.1] (Var4), [0,0.3] (Var5), and [0,0.5] (Var6).

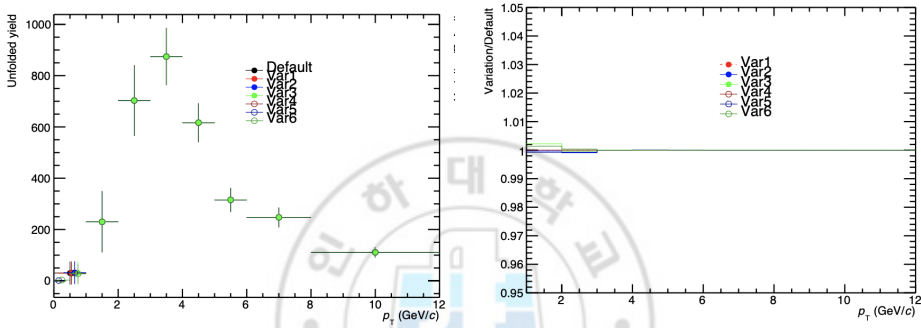


Figure 4.26: Left: the prompt Ξ_c^0 cross section for variations of p_T binning and range. Right: the ratios of the Ξ_c^0 spectra with variations to the standard selection for variations of p_T binning and range.

Fig. 4.26, default binning was fine to describe the low p_T region, and no systematic uncertainty is assigned for p_T binning and range.

Iteration and Singular value decomposition (SVD) method

The number of iterations was 3 in the Bayesian unfolding, which is the default method of this analysis. The number of iterations varied from 2 to 6 to study the uncertainties of the unfolding fluctuations.

The unfolding method was varied to study the uncertainties of the unfolding fluctuations. The SVD method [195] was used to check the uncertainties of the unfolding procedure. Unfolded spectra from the SVD unfolding were sensitive to the d vector. The regularization number was determined as 4 and 5, following the d vector, which is close to 1.

The systematic uncertainties were assigned as a function of p_T , followed by the RMS value.

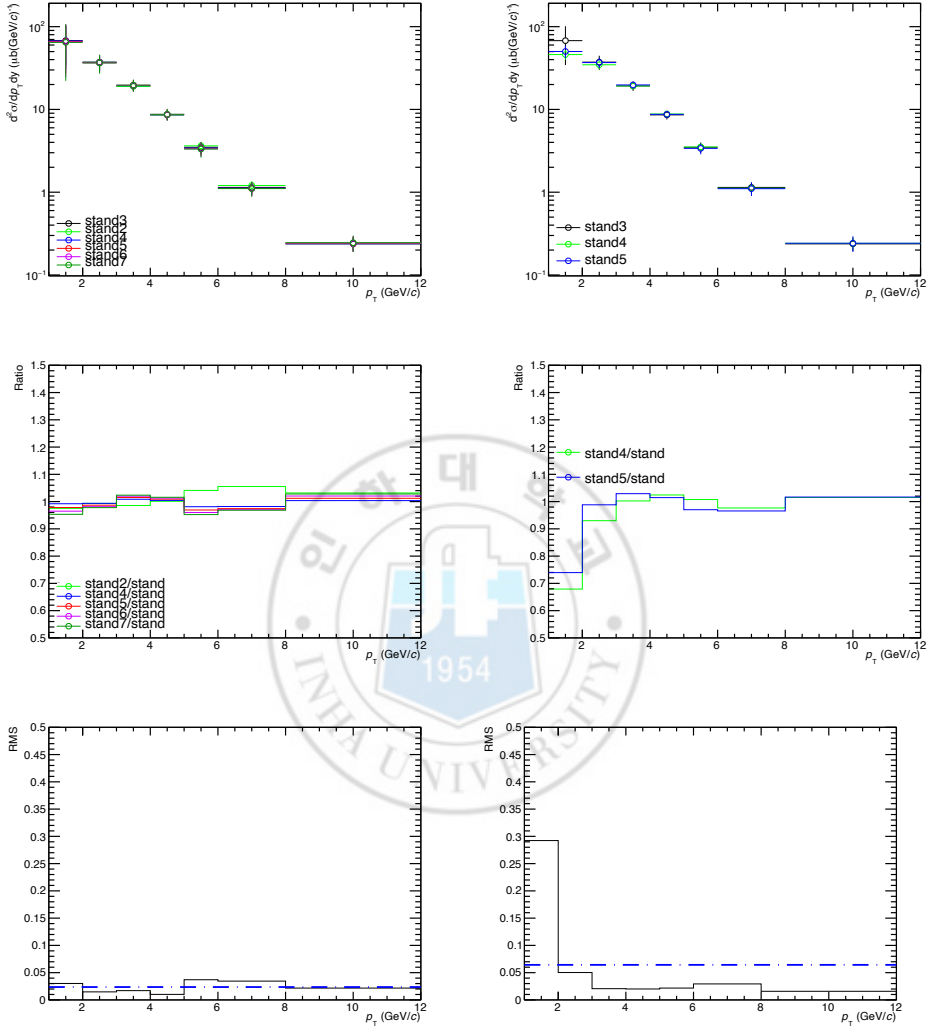


Figure 4.27: Top left (right): the prompt Ξ_c^0 cross section for variations of the number of iterations (the unfolding method). Middle left (right): the ratios of the Ξ_c^0 spectra with variations to the standard selection for variations of the number of iterations (the unfolding method). Bottom left (right): the RMS deviations from 1 for variations of the number of iterations (the unfolding method).

4.4.8 Prompt fraction

The systematic uncertainty on the extraction of prompt fraction (see Sec. 4.3.2) was estimated by varying the FONLL parameters, such as b-quark mass, factorization, and the renormalization scales, and fragmentation functions, and the variation of the ratio of Ξ_c^0 and Λ_c^+ , as maximal value (x2) and minimal value (x0.05) of the ratio, for the conservative calculation.

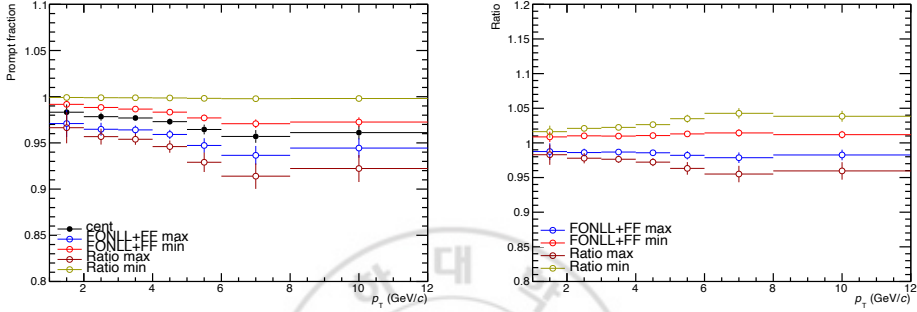


Figure 4.28: Left: the prompt fraction for each variation. Right: the ratios of prompt fractions for Ξ_c^0 .

The systematic uncertainties were assigned as a function of p_T , depending on the variation to the upper and lower direction of the prompt fraction.

4.4.9 ITS-TPC matching efficiency

The systematic uncertainty on the tracking efficiency was estimated by comparing the probability of prolonging a track from the TPC to the ITS ("matching efficiency") in data and simulation. Only the electron track was reconstructed using information from both ITS and TPC detectors. The assumption that the detector response of electron and pion are similar was needed.

Systematic uncertainties of ITS-TPC matching efficiency for electron was propagated to the systematic uncertainties of $e\Xi$ for each p_T bin using Eq. 4.16.

$$\text{Syst}_{\text{ITS-TPCmatching}}^{e\Xi} = \frac{\sum_{i=1}^N (\text{Syst}_i^{e\leftarrow\Xi_c^0} * \text{BC}_i)}{\sum_{i=1}^N \text{BC}_i} \quad (4.16)$$

where BC_i is the Bin Contents of systematic uncertainties of ITS-TPC matching efficiency for electron.

p_T bins	systematic value
1.0 - 2.0	0.018
2.0 - 4.0	0.027
4.0 - 6.0	0.023
6.0 - 8.0	0.024
8.0 - 9.0	0.030

Table 4.11: Systematic uncertainties of ITS-TPC matching efficiency for electron taken from DPG [48] in pp collisions at $\sqrt{s} = 13$ TeV.

$e\Xi$							
p_T bins (GeV/c)	1-2	2-3	3-4	4-5	5-6	6-8	8-12
Systematic	2.52%	2.50%	2.50%	2.44%	2.44%	2.44%	2.50%

Table 4.12: The systematic uncertainty of $p_T^{e\Xi}$ propagated from p_T^e at $\sqrt{s} = 13$ TeV.

Systematic uncertainties of ITS-TPC matching efficiency for $e\Xi$ pair are shown in Table 4.12. Propagating the systematic of $p_T^{e\Xi}$ to the systematic of $p_T^{\Xi^0}$, similarly as before. However, due to the missing neutrino kinematics, the response matrix was considered to connect the $p_T^{e\Xi}$ with $p_T^{\Xi^0}$, which was shown in Fig. 4.12. To get the $p_T^{e\Xi}$ distributions in different $p_T^{\Xi^0}$ intervals, the response matrix was needed to project to the x axis. Systematic uncertainties of ITS-TPC matching efficiency for $e\Xi$ pair were propagated to the systematic uncertainties of $e \leftarrow \Xi_c^0$ for each p_T bin.

$$\text{Syst}_{\text{ITS-TPCmatching}}^{\Xi_c^0} = \frac{\sum_{i=1}^N (\text{Syst}_i^{e\Xi \leftarrow \Xi_c^0} * \text{BC}_i)}{\sum_{i=1}^N \text{BC}_i} \quad (4.17)$$

where BC_i is the Bin Contents of systematic uncertainties of the systematic uncertainty of $p_T^{e\Xi}$.

The systematic uncertainties were assigned as 2.5% for all p_T bins, followed by Table 4.13.

4.4.10 Rapidity range

The rapidity range of Ξ_c^0 in the generation level of MC was varied to study the uncertainties on the Ξ reconstruction from the V0 and cascade.

In the left panel of Fig. 4.29, the red line shows the cross section from the cut

	Ξ_c^0						
p_T bins (GeV/c)	1-2	2-3	3-4	4-5	5-6	6-8	8-12
Systematic uncertainty	2.51%	2.51%	2.50%	2.50%	2.44%	2.45%	2.50%

Table 4.13: The systematic uncertainty of $p_T^{\Xi_c^0}$ propagated from $p_T^{e\Xi} \sqrt{s} = 13$ TeV.

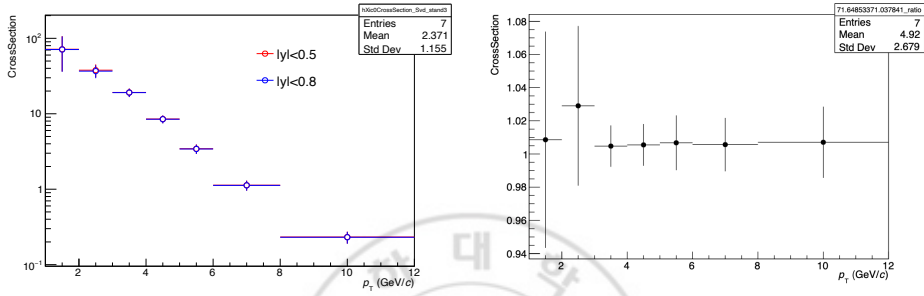


Figure 4.29: Left: the prompt Ξ_c^0 cross section for rapidity range variations. Right: the ratio of the Ξ_c^0 spectra with variations to the standard selection.

$|y| < 0.5$, and the blue line shows the cross section from the cut $|y| < 0.8$. And in the right panel of Fig. 4.29, the ratio of cross sections from each rapidity range.

The systematic uncertainties were assigned as 1% for all p_T bins, followed by the ratio of cross sections.

4.4.11 Summary of systematic uncertainties

The overall systematic uncertainty was calculated by summing the different contributions in quadrature, which were assumed to be uncorrelated.

p_T (GeV/c)	1–2	2–3	3–4	4–5	5–6	6–8	8–12
ITS-TPC matching	2.5%						
Electron reconstruction	1%						
$\Xi^\pm$ -daughter reconstruction	7%	8%	1%	1%	2%	1%	1%
Electron identification	12%	5%	3%	3%	2%	3%	9%
$\Xi^\pm$ topological selection	6%	8%	5%	3%	2%	5%	5%
e Ξ -pair mass	13%	7%	5%	10%	14%	9%	9%
opening angle	20%	4%	4%	3%	1%	0%	0%
Bayesian-unfolding iterations	3%	1%	1%	4%	3%	2%	2%
Unfolding method	29%	5%	2%	2%	2%	3%	2%
Over subtraction	1%	1%	2%	2%	5%	6%	6%
Generated p_T shape	1%	1%	1%	1%	1%	1%	3%
Rapidity interval	1%						
Feed-down subtraction (λ_c^+ variation)	$+1\%$ -1%	$+1\%$ -1%	$+1\%$ -1%	$+1\%$ -1%	$+1\%$ -2%	$+1\%$ -2%	$+1\%$ -2%
Feed-down subtraction (Ξ_c^0 variation)	$+2\%$ -2%	$+2\%$ -2%	$+2\%$ -2%	$+3\%$ -3%	$+4\%$ -3%	$+5\%$ -4%	$+4\%$ -4%
Branching ratio	[196]						
Luminosity	[49, 197]						

Table 4.14: Contributions to the systematic uncertainty on the Ξ_c^0 cross section.

4.5 Results

4.5.1 Branching-fraction ratio

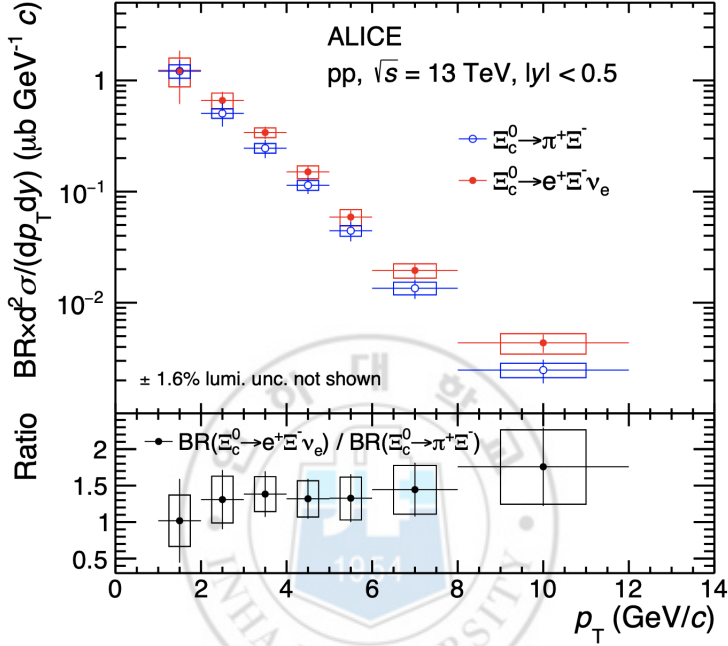


Figure 4.30: Upper panel: BR x cross sections of prompt $\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e$ (full red circles) and $\Xi_c^0 \rightarrow \Xi^- \pi^+$ (open blue circles) as a function of p_T . Lower panel: ratio of BR x cross sections of prompt $\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e$ and $\Xi_c^0 \rightarrow \Xi^- \pi^+$ as a function of p_T . Figures from [37].

The production cross section of prompt Ξ_c^0 baryons without BR correction in both the semileptonic ($\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e$) and hadronic decay ($\Xi_c^0 \rightarrow \Xi^- \pi^+$) channels are shown in the upper panel of Fig. 4.30. The error bars indicate the statistical uncertainties, and the boxes indicate the statistical and systematic uncertainties. The Ξ_c^0 measurements in the two decay channels are compatible within statistical and systematic uncertainties.

The lower panel of Fig. 4.30 shows the p_T -dependent ratio of the two Ξ_c^0 measurements. The average of uncertainty for the p_T -dependent ratio of the two Ξ_c^0 measurements was calculated as the sum in quadrature of the relative statistical uncertainties and the weight of the p_T -uncorrelated sources of the systematic un-

certainties [198]. The systematic uncertainty of unfolding, for the $\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e$, was considered as p_T uncorrelated uncertainties, and the other uncertainties were considered as fully p_T correlated uncertainties. The systematic uncertainty of the production ratio was calculated with the quadrature sum of p_T -correlated and uncorrelated systematic uncertainties. The weighted average of the production ratio was considered to calculate the branching fraction ratio:

$$\text{BR}(\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e) / \text{BR}(\Xi_c^0 \rightarrow \Xi^- \pi^+) = 1.38 \pm 0.14 (\text{stat}) \pm 0.22 (\text{syst}). \quad (4.18)$$

with an improved precision by a factor of three with respect to the global average value reported by the PDG [49]. The branching fraction ratio which was one by the Belle collaboration [199] shows a compatible value with the one measured by ALICE collaboration.

4.5.2 Production cross section

The measurement of the prompt Ξ_c^0 cross section at $\sqrt{s} = 13 \text{ TeV}$ using the two decay channels considering each branching ratio is shown in fig. 4.31. The cross sections of the prompt Ξ_c^0 for the two decay channels were averaged to obtain a more precise measurements of the p_T -differential Ξ_c^0 production cross section. The average of the cross sections is calculated considering as weights the inverse square of the relative statistical uncertainties and the p_T -uncorrelated systematic uncertainties [198]. The cross sections of Ξ_c^0 in the two decay channels and the average cross section of Ξ_c^0 are consistent within the uncertainty of the measured p_T interval. The error bars indicate statistical uncertainties, and empty boxes indicate systematic uncertainties. The systematic uncertainties of the branching ratio are shown as shaded boxes, and the integrated luminosity uncertainties are shown as a value in Fig. 4.31.

The cross section of prompt Ξ_c^0 in pp collisions at $\sqrt{s} = 13 \text{ TeV}$, measured in the rapidity interval $|y| < 0.5$ and p_T range $1 < p_T < 12 \text{ GeV}/c$, is shown in the Fig. 4.33. The prompt Ξ_c^0 cross section is compared with the prompt Ξ_c^+ cross section as shown in Fig. 4.31, which is the isospin partner. The measurements of prompt Ξ_c^0 and Ξ_c^+ are compatible within uncertainties.

The prompt Ξ_c^0 production at $\sqrt{s} = 13 \text{ TeV}$ is compared with the inclusive Ξ_c^0 production at midrapidity in pp collisions at $\sqrt{s} = 7 \text{ TeV}$, and the prompt Ξ_c^0 production at midrapidity in pp collisions at $\sqrt{s} = 5.02 \text{ TeV}$ [39]. The error bars indicate statistical uncertainties, and empty boxes indicate systematic uncertainties. The systematic uncertainties of the branching ratio are shown as shaded boxes, and the integrated luminosity uncertainties are shown as a value in Fig. 4.33. The uncertainty on the branching ratio of the prompt Ξ_c^0 at $\sqrt{s} = 7 \text{ TeV}$ and $\sqrt{s} = 5.02 \text{ TeV}$ decreases because it is constrained by the new measurement of a combination of two different

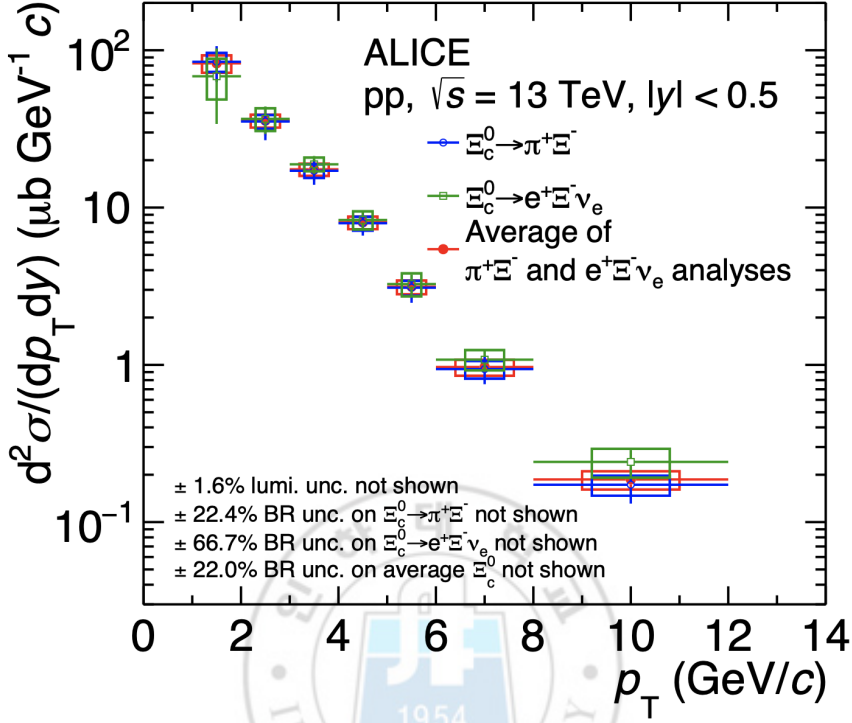


Figure 4.31: Measurement of prompt cross sections for $\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e$ (empty green squares) and $\Xi_c^0 \rightarrow \Xi^- \pi^+$ (open blue circles), and the measurement of average of the cross section (full red markers) as a function of the p_T . The uncertainty on integrated luminosity is not included. Figure from [37].

decay channels ($\Xi_c^0 \rightarrow e^+ \Xi^- \nu_e$ and $\Xi_c^0 \rightarrow \pi^+ \Xi^-$) at $\sqrt{s} = 13$ TeV. The prompt fraction in the Ξ_c^0 is close to unity, as shown in Sec. 4.3.2. Therefore, the comparison of the inclusive Ξ_c^0 cross section measured at $\sqrt{s} = 7$ TeV with the prompt ones at $\sqrt{s} = 5$ and 13 TeV can provide a meaningful message of the $\sqrt{s}$ dependence for the production. The $\sqrt{s}$ dependence of the production indicates a hardening of the p_T -differential cross section with increasing collision energy. The trend of $\sqrt{s}$ dependence is consistent with the observation for the D-meson and Λ_c^+ at $\sqrt{s} = 5.02$, 7 and 13 TeV [120, 200, 156, 25, 201], and with the pQCD calculations [114].

The visible cross section of Ξ_c^0 at midrapidity in pp collisions at $\sqrt{s} = 13$ TeV was calculated in the measured p_T interval by integrating the measured Ξ_c^0 production. The branching ratio uncertainty was considered as a systematic uncertainty.

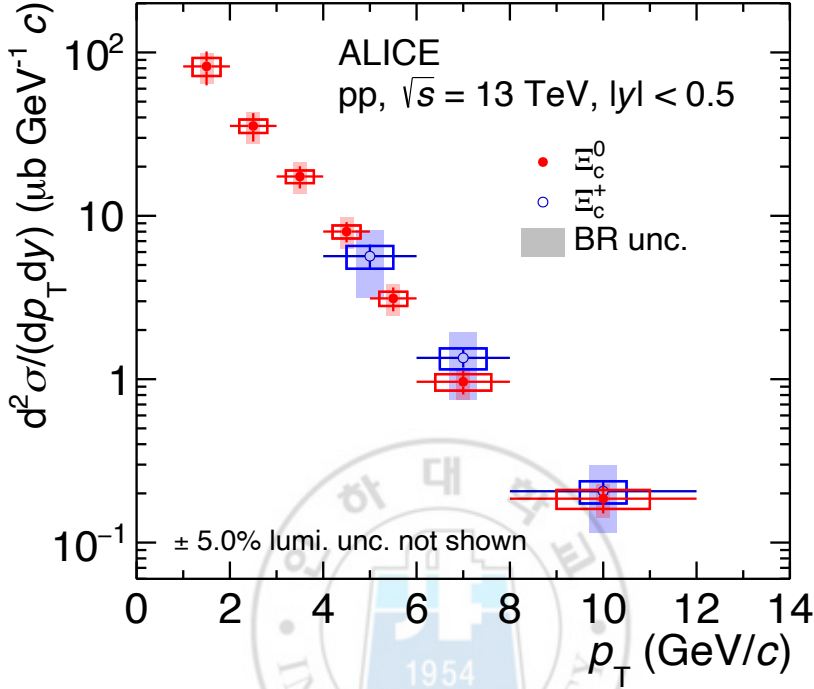


Figure 4.32: The measurement of prompt cross sections for Ξ_c^0 (full circles) and Ξ_c^+ (open circles) baryons as a function of p_T in pp collisions at $\sqrt{s} = 13$ TeV. The statistical and systematic uncertainties are represented by the error bars and empty boxes. Shaded boxes represent the systematic uncertainties of the branching ratio (BR). The boxes do not contain the integrated luminosity uncertainty.

The p_T -integrated cross sections of Ξ_c^0 in the measured p_T interval is:

$$\left. \frac{d\sigma_{pp, 13 \text{ TeV}}^{\Xi_c^0}}{dy} \right|_{|y| < 0.5}^{(1 < p_T < 12 \text{ GeV}/c)} = 148 \pm 21 \text{ (stat.)} \pm 35 \text{ (syst.)} \pm 7 \text{ (lumi.) } \mu\text{b.} \quad (4.19)$$

The p_T -extrapolated Ξ_c^0 cross section at midrapidity in pp collisions at $\sqrt{s} = 13$ TeV was extracted by extrapolating the measured cross section to the full p_T range. The Catania model [27], was used to calculate the extrapolation factor, which is $1.36^{+0.05}_{-0.15}$, which could better explain the measured p_T shape. The systematic uncertainty was estimated considering calculations [26, 46, 202] that describe the p_T

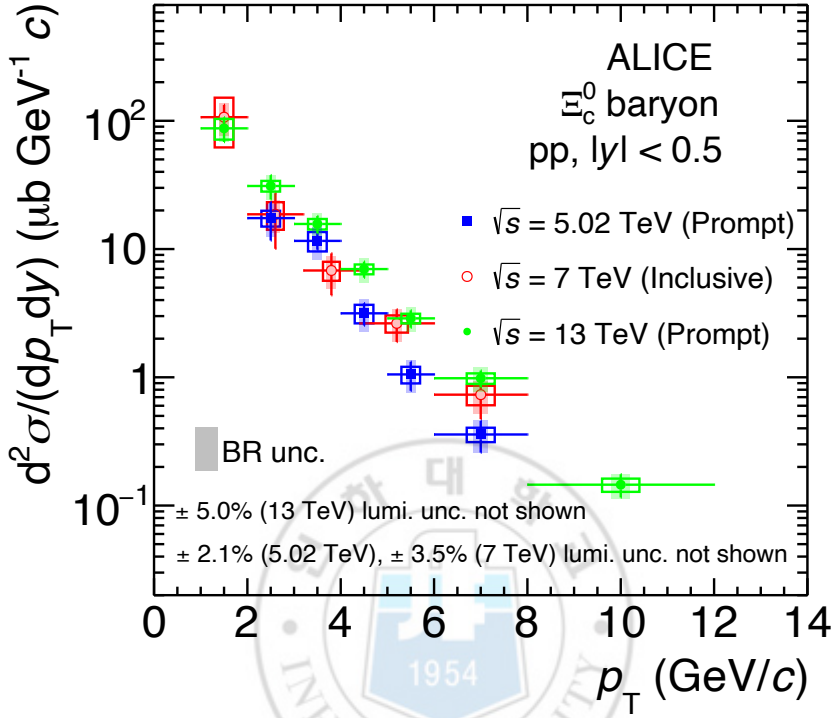


Figure 4.33: p_T -differential production cross sections of prompt Ξ_c^0 baryons in pp collisions at $\sqrt{s} = 5.02$ TeV and 13 TeV and of inclusive Ξ_c^0 baryons in pp collisions at $\sqrt{s} = 7$ TeV [38]. Figure from [39].

shape of the measured cross section. The p_T -extrapolated cross section of Ξ_c^0 is:

$$\left. \frac{d\sigma_{pp, 13 \text{ TeV}}^{\Xi_c^0}}{dy} \right|_{|y| < 0.5} = 201 \pm 28 \text{ (stat.)} \pm 48 \text{ (syst.)} \pm 10 \text{ (lumi.)} {}^{+10}_{-21} \text{ (extrap.) } \mu\text{b}. \quad (4.20)$$

The production of Ξ_c^0 baryons shows a sizable contribution to the $c\bar{c}$ production cross section at midrapidity per unit rapidity in pp collisions at the LHC [203]. The p_T -integrated cross section of Ξ_c^0 is used to derive the ratio of the Ξ_c^0 and D^0 meson, which is measured at the same collision energy [120].

4.5.3 Baryon-to-meson and baryon-to-baryon ratio

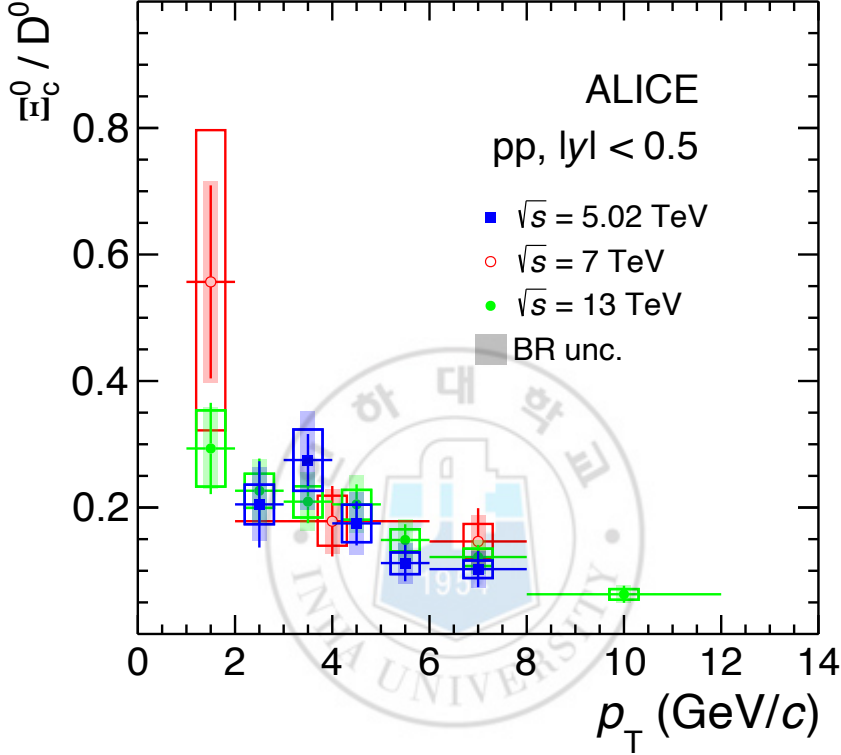


Figure 4.34: Ξ_c^0/D^0 ratio measured in pp collisions at $\sqrt{s} = 13$ TeV (green), compared with the measurements at $\sqrt{s} = 7$ TeV [38] (red) and $\sqrt{s} = 5.02$ TeV [39] blue. The uncertainty of the branching ratio (BR) of D^0 is included in the point-by-point systematic uncertainties in the empty boxes, while the BR uncertainties on the Ξ_c^0 are shown as shaded boxes. Figure from [39].

The ratio of the p_T -differential Ξ_c^0 and D^0 yield at midrapidity measured in pp collisions at $\sqrt{s} = 13$ TeV is shown in Fig. 4.34, and compared with the same baryon-to-meson ratio at midrapidity measured in pp collisions at $\sqrt{s} = 7$ [38] and 5.02 TeV [39]. The prompt D^0 cross section was published in Ref. [204], which is used to evaluate the charm baryon-to-meson ratio, matching the p_T bin. The systematic uncertainties on D^0 meson were propagated considering the raw yield extraction uncertainty as fully uncorrelated and other sources of uncertainties as fully

correlated in the p_T intervals when rebinning the D^0 cross section. The systematic uncertainty on the Ξ_c^0/D^0 ratio was calculated with the assumption that all the uncertainties of the Ξ_c^0 and D^0 cross sections as uncorrelated. However, the ratio considers tracking and feed-down systematic uncertainties partially canceled. The luminosity uncertainty is fully canceled in the Ξ_c^0 and D^0 ratio. The Ξ_c^0/D^0 ratios at the three center-of-mass energies are compatible with uncertainties, while the prompt cross shows the $\sqrt{s}$ dependence. The Ξ_c^0/D^0 ratio shows the values about 0.35 at low p_T , and it decreases with increasing p_T , converging a value of about 0.1 for $p_T > 8$ GeV/c. The Ξ_c^+/D^0 ratio is also shown in Fig. 4.35, and the ratio is compatible with the Ξ_c^0/D^0 within uncertainties.

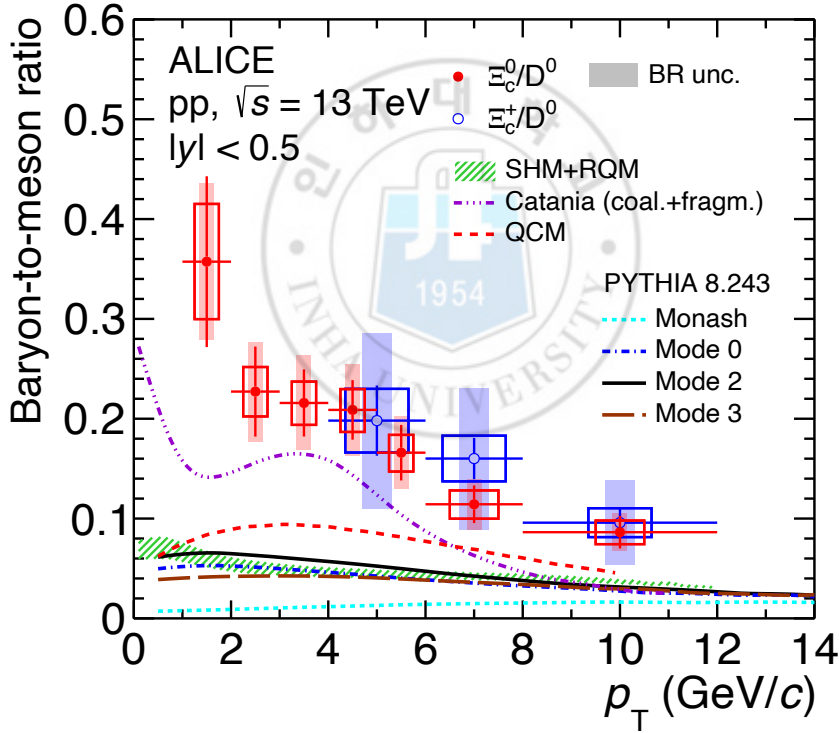


Figure 4.35: Ξ_c^0/D^0 (red) and Ξ_c^+/D^0 (blue) ratios as a function of p_T in pp collisions at $\sqrt{s} = 13$ TeV. While the branching ratio (BR) uncertainties on the Ξ_c^0 are presented as shaded boxes, the BR uncertainty of the D^0 is included in the point-by-point systematic uncertainties in the empty boxes. Figure from [37].

The measured Ξ_c^0/D^0 ratio is also compared with the model calculations. The measured Ξ_c^0/D^0 ratio is compared in Fig. 4.35 with the calculations from the different tunes of the PYTHIA 8 event generator. The Monash tune significantly underestimates the data by a factor of about 23–43 in the low p_T region and by a factor of about 5 in the highest p_T interval at $\sqrt{s} = 13$ TeV in pp collisions, as also confirmed for the Λ_c^+/D^0 ratio [191]. All three CR modes yield a similar magnitude and shape of the Ξ_c^0/D^0 ratio. Despite expecting a more considerable Ξ_c^0/D^0 ratio with respect to the Monash tune, the CR modes still underestimate the measured Ξ_c^0/D^0 ratio by a factor of about 4–9 at low p_T .

The SHM calculation [46] describes the baryon formation based on the thermal weights considering the additional excited charm-baryon states not yet observed in pp collisions. The feed-down from the excited charm-baryon states can increase the chance of composing the charm-baryon. The additional charm-baryon states are counted based on the relativistic quark model (RQM) [144] and lattice QCD calculations [205]. By varying the branching ratios of excited charm-baryon states decaying to the ground state $\Xi_c^{0,+}c$, the uncertainty band of the model is evaluated. The combined model of SHM and RQM does not describe both measurements, the Ξ_c^0/D^0 , and Ξ_c^+/D^0 .

The production cross section of the Ξ_c^0 baryon is compared with a model considering the coalescence process, QCM [202], charm quarks are combined with same velocity hadrons. It also underestimates the measurement of the Ξ_c^+/D^0 ratio by a factor 2–3 for $p_T < 4$ GeV/c.

The Catania model [27] considered hadronization of charm-quark via vacuum fragmentation and coalescence in pp collisions (see Sec. 2.2.3). In this model, the coalescence is described with the Wigner formalism, according to which the charm coalescence produces the p_T spectrum of a hadron. The Catania model describes the Ξ_c^0/D^0 ratio considering the uncertainties in the whole p_T interval of the measurement.

The charmed baryon-to-baryon ratio, Ξ_c^0/Λ_c^+ and $\Xi_c^{0,+}/\Sigma_c^{0,+,++}$ [156] ratios are shown in Fig. 4.36. The tracking, feed-down, and luminosity uncertainties were considered correlated, and the rest of the uncertainties source were considered uncorrelated. The evaluation of the $\Xi_c^{0,+}/\Sigma_c^{0,+,++}$ was done by the sum of Ξ_c^0 and Ξ_c^+ for $p_T > 4$ GeV/c and scaling by a factor of two in the interval $2 < p_T < 4$ GeV/c. The Ξ_c^0/Λ_c^+ ratio is around 0.5, and $\Xi_c^{0,+}/\Sigma_c^{0,+,++}$ ratio is around 2. There is no strong p_T dependence within the current uncertainties. All the PYTHIA 8 tunes, QCM, and SHM+RQM models calculation can not describe the charmed baryon-to-baryon ratio, except for the $\Xi_c^{0,+}/\Sigma_c^{0,+,++}$ ratio. The $\Xi_c^{0,+}/\Sigma_c^{0,+,++}$ ratio is compatible with the Monash tune due to a similar amount underestimation on the $\Xi_c^{0,+}$ and $\Sigma_c^{0,+,++}$ cross sections [38, 156]. The PYTHIA 8 tunes with CR and the SHM+RQM calcu-

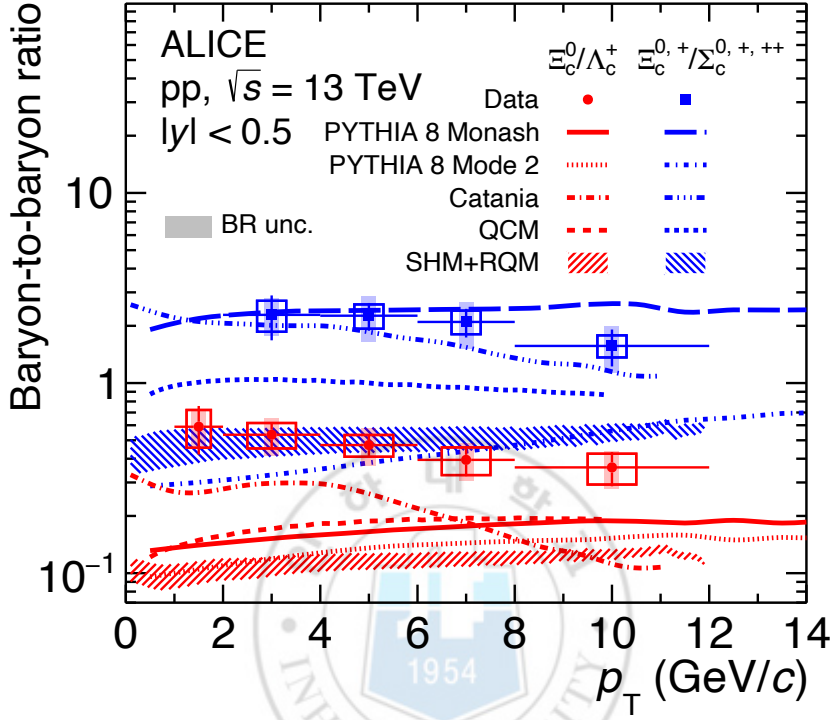


Figure 4.36: Ξ_c^0/Λ_c^+ and $\Xi_c^{0,+}/\Sigma_c^{0,+,++}$ ratio as a function of p_T at $\sqrt{s} = 13$ TeV. The statistical and systematic uncertainties are represented by the error bars and empty boxes. Shaded boxes represent the systematic uncertainties of the branching ratio (BR). The measurements and model calculations are compared. Figure from [37].

lation also underestimate the measurements of Ξ_c^0/Λ_c^+ and $\Xi_c^{0,+}/\Sigma_c^{0,+,++}$ ratios. The QCM model shows an almost flat value at unity, significantly underestimating both measured ratios. The Catania model describes the data within the uncertainties, as did in Ξ_c^0/D^0 ratio.

Chapter 5

Ξ_{cc}^{++} performance in pp collisions at $\sqrt{s} = 14\text{TeV}$ with ALICE 3

All four main LHC experiments have been participating in the ultra-relativistic heavy-ion physics program, showing the advanced properties of the QGP. Despite the ambitious scientific program for the upcoming Run 3 and 4, crucial questions will still need to be answered with the current detector concepts. The aims of ALICE 3 are to understand the rich phenomenology of QCD matter, connection parton energy loss, collective flow, hadronization, and electromagnetic radiation in a unified description.

The physics goal and innovative detector concept will be introduced in Sec. 5.1, and the analysis framework, which will handle all of ALICE 3 data analysis, will be introduced in Sec. 5.2. One of the physics goals of ALICE 3 is searching for multi-charm baryons. In Sec. 5.3.3, the results of the performance study for $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ will be discussed. The performance study is done using the machine learning technique introduced in Sec. 5.3.1.

5.1 ALICE 3

ALICE experiment has already led to a deep understanding of the properties of the QGP in Run 1 and 2 at LHC. Despite the considerable progress made so far (see Sec. 1.4), several fundamental questions still remain and need to be studied. The dynamical mechanisms, partons lose momentum, hydrodynamic, and thermalize in dense QCD matter still need to be discovered. In particular, the hadronization of the heavy-flavor baryons still needs to be completed, as shown in Sec. 4.5.

ALICE 3 proposes an innovative detector concept to manage the uncovered physics topics. It has a high-speed silicon tracker with high resolution and minimum mass placed close to the collision point. The first layers are placed inside the beam pipe and provide a very wide rapidity acceptance, complemented with hadron, electron, muon, and photon identification performance over a near the 0 to high p_T range.

5.1.1 Physics goal

ALICE experiment aims to specify the properties of strongly interacting matter and determine how they arise from the underlying interactions described by quantum chromodynamics. Despite the expected progress in the current decade, it will not be possible to understand the undiscovered area with the current or currently planned detectors. Therefore, in 2035, ALICE 3 will operate with a novel experimental approach to study the fundamental physics of QGP. A few key examples are given below.

High-precision beauty measurements

Beauty quarks may be more appropriate for pQCD calculation to describe the processes involving the QCD medium due to their significantly larger masses. The precision measurements of the cross section and flow coefficients of beauty hadrons are necessary down to the momenta typical of diffusion phenomena to establish a firm connection between parton transport, collective flow, and hadronization. Also, beauty hadrons can provide a much cleaner field of study with much fewer feed-down contributions than charm hadrons.

$D\bar{D}$ correlations

The correlations between fully reconstructed charm hadron pairs over a wide rapidity range are required to distinguish between the different regimes of in-medium energy loss and establish the onset of charm isotropization.

Multi-charm baryons, P-wave quarkonia, exotic hadrons

The measurement of multi-heavy-flavor hadrons, P-wave quarkonia, and exotic states will offer unprecedented sensitivity to studying the formation of hadrons from the deconfined QGP. Comparisons between S-wave and P-wave states for charmonium and bottomonium provide a more detailed characterization of dissipative effects. The production of exotic states inside a deconfined medium provides information on the formation and interaction of exotic states, the characterization of binding potential, and the details of their hadronization mechanisms.

Azimuthal asymmetry of electromagnetic radiation

Studying the azimuthal asymmetries of different particle species is a potent tool to access the shear and bulk viscosities by correlating final-state anisotropies to the geometry of the initial medium. Moreover, electromagnetic probes can offer information on full-time evolution and their temperature dependence since they do not have strong interaction with the medium. Therefore, azimuthal asymmetry of the dileptons as a function of transverse momentum and mass provides a direct connection to the emission time.

Chiral symmetry restoration

The a_1 axial meson, which is the chiral partner of the ρ meson, does not couple to the dilepton channel in the vacuum. However, as chiral symmetry is restored in the deconfinement medium, a_1 axial meson starts to mix with the ρ meson, leading to a significant modification of the dielectron mass spectrum.

Collectivity in small systems

The finding of collective effects in small collision systems has been one of the most exceptional LHC results. The investigation of the source of collective effects in small collision systems will be one of the major goals in the high-multiplicity pp program in ALICE 3.

High-statistics hadronic physics

High statistics will allow access to rare observables, such as the production of nuclei, hypernuclei, and super-nuclei, the nuclei in which a charmed baryon replaces a nucleon. Wide acceptance, with excellent reconstruction at low transverse momenta, will extend the investigations of hadron-hadron potentials via two-particle

correlations to the charm sector, providing a structure of the newly-discovered charmed exotic states. Moreover, the detector concept in ALICE 3 allows for exploring more complex vector meson decays in ultra-peripheral collisions and the reconstruction of photon conversion at ultra-low transverse momenta.

5.1.2 Detector setup in ALICE 3

ALICE 3 is proposed as a novel detector with a high read-out rate, excellent pointing resolution, and outstanding tracking and particle identification over a large acceptance using advanced silicon detectors. This combination provides unique access to thermal dielectron production and heavy flavor probes of the quark-gluon plasma.

The first tracking layer will be positioned close to the interaction point for excellent pointing resolution. The vertex detector will be retractable to provide the larger aperture required for the beams at injection energy. The vertex layers are made from wafer-scale CMOS Active Pixel Sensors thinned to $30\mu\text{m}$ and bent into cylinders to minimize the material. An outer tracker with barrel and endcap layers provides a relative momentum resolution of 1-2% over a large acceptance.

Particle identification will be done by a time-of-flight detector (TOF) and a ring-imaging Cherenkov detector (RICH) covering a wide momentum range, relying on novel silicon timing and photon sensors. An electromagnetic calorimeter and a muon identifier will provide photon detection and lepton identification at higher momentum. A forward conversion tracker will measure photons at very low p_T in forward rapidity.

5.2 Analysis framework in ALICE 3

ALICE 3 detector plans to collect the extensive Pb–Pb statistics provided by the LHC up to integrated luminosities per one month as $\mathcal{L} = 242\text{ pb}^{-1}$, corresponding to an interaction rate of about 93 kHz [24]. The computing model is designed to reduce the data volume to proceed with the readout from the detector as early as possible during the data flow, minimizing the cost, data processing, and storage. And data analysis and MC production is done with Analysis Object Data (AOD) files. All data processing is handled by ALICE Online-Offline (O^2) analysis framework [206].

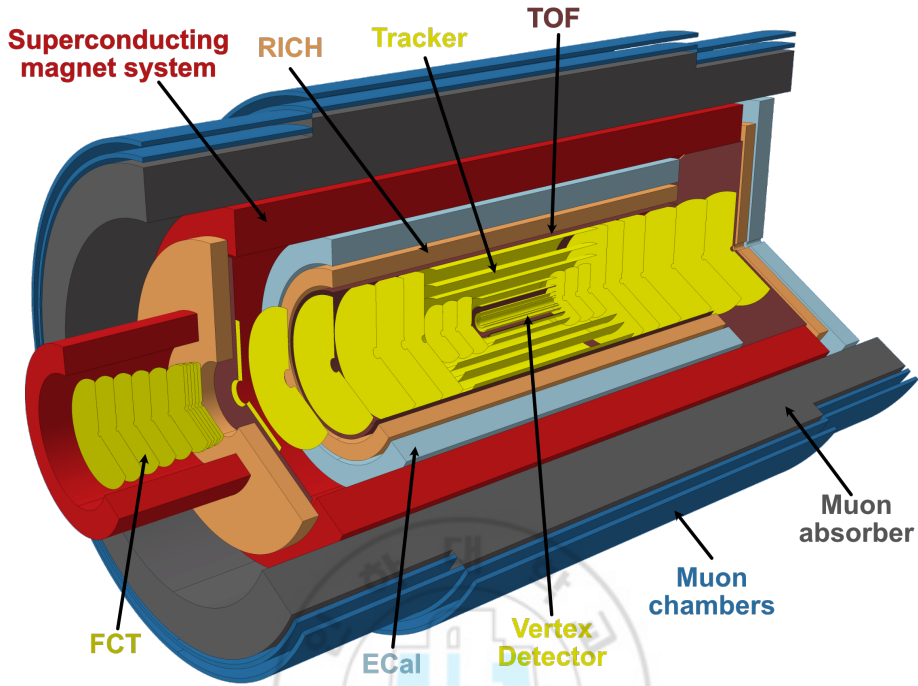


Figure 5.1: A silicon tracker made of cylinders and disks is used in the ALICE 3 detector to reconstruct tracks in the magnetic field produced by a superconducting magnet system. The vertex tracker is included within the beam pipe. For particle identification, a time-of-flight detector, RICH detector, photon detector, and muon system are used. Figure from [24].

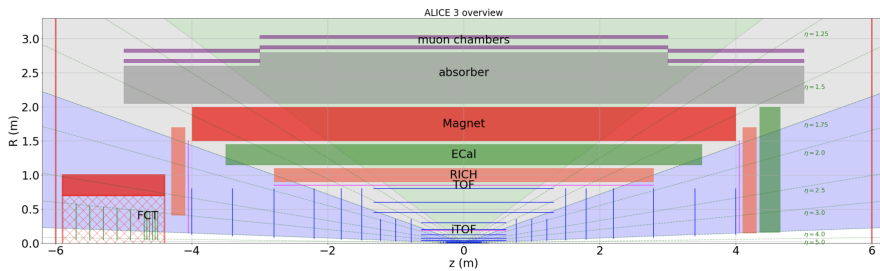


Figure 5.2: Longitudinal cross section of the ALICE 3 detector concept. Figure from [24].

5.2.1 Basic features of the (O^2) analysis framework

The (O^2) software framework is designed from the start to combine all the computing functionalities needed in a HEP experiment: detector read-out, event building, data recording, detector calibration, data reconstruction, physics simulation, and analysis [207]. In the following section, analysis workflow in (O^2) will be focused on, including all the utilities needed to perform the complete analysis chain, from processing the reconstructed events stored in AOD files to extracting the histograms used in the final analysis.

5.2.2 Heavy-flavour reconstruction, skimming, and selection

A detailed sketch of the Heavy-Flavor (HF) analysis workflow in (O^2) is presented in the diagram in Fig. 5.3. Details of each step are introduced in the following subsections.

Track pre-selection

The track tables are accessed and filtered by quality and kinematic selections. Loose selections on the track p_T and DCA are applied to consider HF candidates down to very low p_T , based on the displacement from HF decays and primary vertex.

Track combination and proto-candidate pre-selection

Double and triple tracks of each event are combined to reconstruct hadrons that decay in two or three charged particles (or prongs) by adding the 4-momenta of the considered tracks, and selections of invariant mass or p_T are applied based on hypotheses on the daughter-track masses.

First secondary-vertex reconstruction and candidate building

The decay vertex (or secondary vertex) is reconstructed by minimizing the distance of the closest approach of the selected tracks to the vertex using the Newton-Raphson method. If the secondary vertex is found, the composed tracks are propagated to the decay vertex using a Kalman algorithm.

A first selection at the HF candidate level is applied based on its topologies like candidate p_T , cosine of the pointing angle, and product of prong impact parameters. For the selected candidates, the indices of the decay tracks and decay channel are stored as columns in 'TrackIndexSkim'.

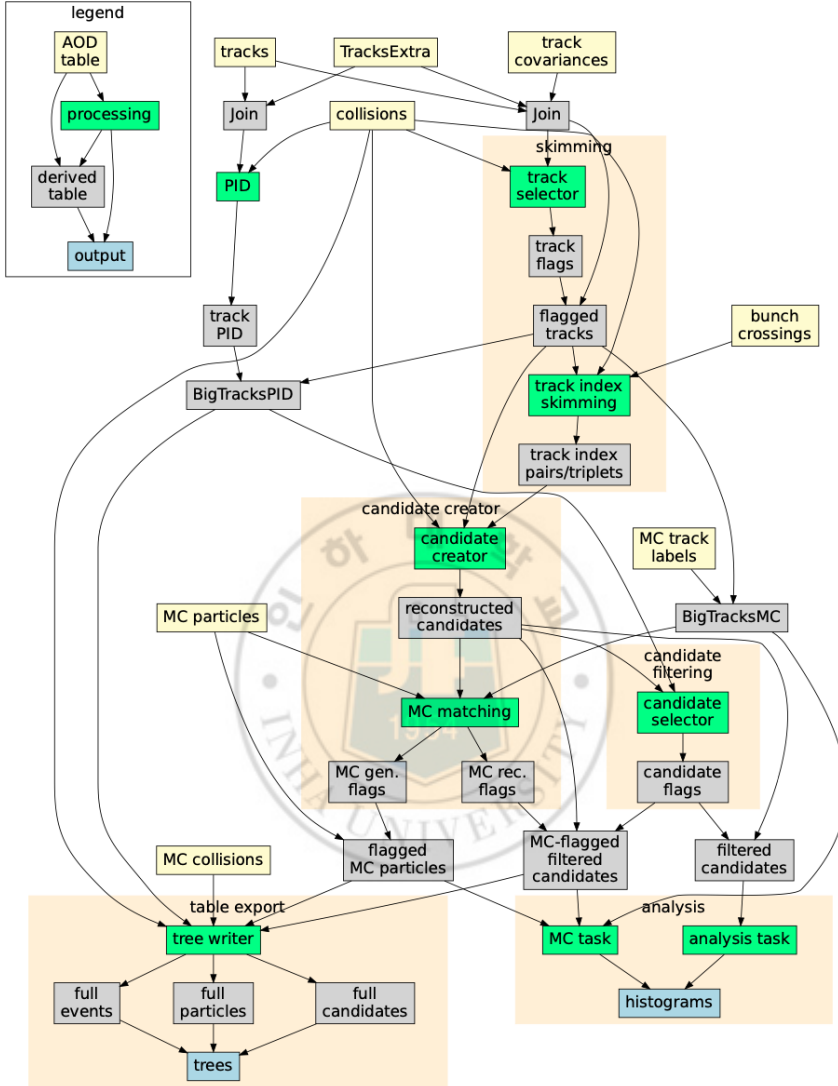


Figure 5.3: Diagram of the HF analysis framework in (O^2). Figure from [40].

Candidate rebuilding and Monte-Carlo matching

The track indices stored for each candidate are used to get track and topology parameters from the AOD track tables. The secondary vertex reconstruction is re-

peated, and additional candidate properties are needed for the signal selection, and cannot be calculated dynamically on-the-fly, are calculated. The list of reconstructed candidate properties is stored in a derived table called 'CandidateSkims'.

For simulated events, reconstructed candidates are matched with their generated partners by MC information of their identity and decay tree. At this stage, derived tables with MC flags are used to estimate the signal efficiencies and optimize the signal and background selections.

Final candidate selection and histogram filling

Analysis level selection criteria based on decay topology and PID information are applied to the reconstructed candidates by the dedicated selector task for each decay channel. The selection flag is joined with the candidate table to filter them.

The final workflow is the user analysis task, where the histograms needed for the analysis are saved in ROOT files. Histograms of properties for generated MC particles and MC-matched candidates are also constructed for MC events.

5.3 $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ reconstruction

5.3.1 Machine learning

Machine learning (ML) is a branch of artificial intelligence (AI) dealing with algorithms (models) that learn to perform a task. ML is typically used to solve classification problems, in which two or more classes of events must be distinguished, and regression problems, in which the algorithm makes predictions based on the input data [41]. The most common binary categorization in high-energy physics is signal and combinatorial background separation. In supervised learning, the algorithm is trained to recognize the various classes in labeled data (training sample), also known as training. The model parameter configuration for the best event categorization, referred to as features, is established through training. The topological and PID selections of daughter tracks are among the features used in reconstructing the $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ decays. The algorithm should ideally discover general patterns specific to the various classes to classify unlabelled datasets. The score, a numerical number calculated for each event in accordance with the features, is the output of the model training. The class selection is decided by comparing this value to a fixed threshold.

In the subsequent session of this research, Boosted Decision Tree (BDT) is used to create a binary classification model that separates the XiccDecay signal from the background. The $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ signal sample is generated via

Delphes O2 Monte Carlo simulations, and the combinatorial background sample is generated for MB events by $pK - \pi^+$ sets for the training and testing samples. The model is applied to the data to search for the $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ signal.

5.3.2 Boosted Decision Trees(BDT)

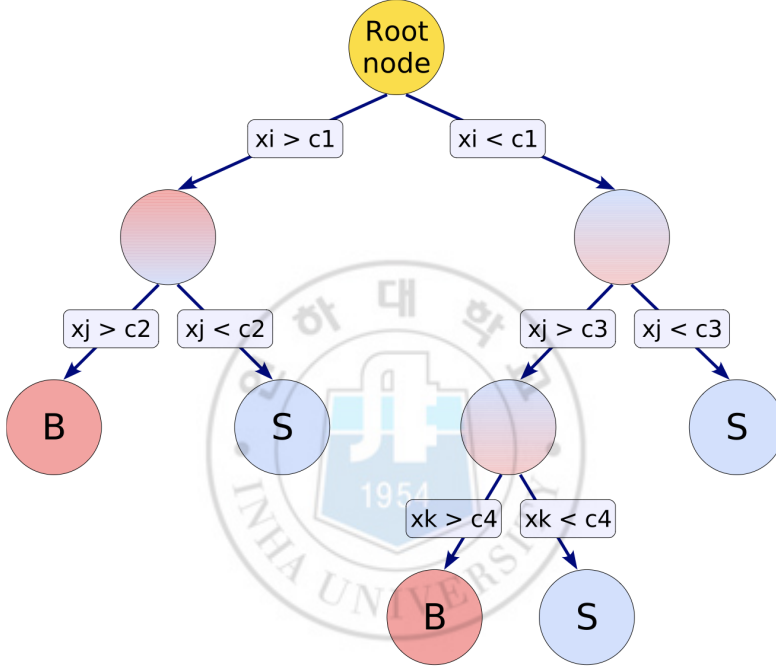


Figure 5.4: Structure of a decision tree built to distinguish signal events (S) from background ones (B). Figure from [41].

The Boosted Decision Trees (BDT), constructed from the Decision Tree (DT) block, is a widely used approach for binary classification [41]. A DT is a sequential model that combines simple categorization to discriminate between two events based on their features. A DT's basic structure is sketched in Fig. 5.4. The decision (test) determines the difference between the signal (S) and the background (B). An initial test is run at the root node from which the tree grows. The threshold value for a specific condition and the numeric feature are compared in the test to distinguish the two classes. The classified event is sent in one of the two directions (left or right) based on the test results, and the procedure is repeated with the other test condition

until it reaches the stopping criterion established a priori by the user.

The DT construction is performed automatically during the training, and a G_{ini} index is commonly used in the node, defined as:

$$G_{\text{ini}} = \left(\sum_{i=1}^N w_i \right) P(1 - P) \quad (5.1)$$

where w_i is the weight of the i -th event, and P is the purity of the signal in the considered node. The criterion for the best separation consists in minimizing the G_{ini} index difference between the parent and children nodes, represents as:

$$c = G_{\text{ini}} - (G_{\text{ini}}^{\text{leftchild}} + G_{\text{ini}}^{\text{rightchild}}) \quad (5.2)$$

The c minimization proceeds sequentially, searching for the features and selection criteria providing the best splitting at each node. In order to accurately classify the training dataset, a huge number of DTs may be generated in this manner. However, there is a strong likelihood that the overall patterns defining the two classes will be missed. Because of this, the depth of DT is typically constrained to prevent overtraining. A small fluctuation may impact the algorithm stability in the training dataset in the small DT scenario because the discrimination of small DT is constrained. In order to create a strong model with good discrimination, additional DTs are utilized, and the final classification is obtained by merging the results of all the DTs. Several DTs are used under the direction of a Random Forest [208], which produces a classification from the (weighted) average of the results from all the DTs in the forest. A feature ranking was conducted based on how frequently each feature was used throughout DT construction. In high-energy physics, feature ranking may represent the critical variables with a precise physical meaning to separate the signal and background events. The creation of hundreds of independent DTs requires such huge computational resources. Moreover, Adding a new DT helps to improve the whole algorithm's discrimination power without information from the previous classifier, which does not improve the classification. For this reason, boosting techniques are required.

In Boosted Decision Trees (BDT) methods, the DTs are built sequentially, attempting to add a new DT that supplements the information of the previous classifier. BDT is based on the event re-weighting to increase the importance of the events which is misclassified by the previous classifier. The BDT model equivalents the ensemble of DTs, and the outcomes are more stable than a single DT to the training-set fluctuations. The BDT provides an output score corresponding to each DT's outcome. The most popular BDT algorithm is XGBoost (XGBoost) [209],

adopted in this analysis. XGBoost stands for “Extreme Gradient Boosting”, where the term “Gradient Boosting” originates from the Ref. [210].

5.3.3 Search for Ξ_{cc}^{++} signal

The production of the multi-charm baryon is suppressed in pp collisions while shows enhancement towards the AA collision, discussed in Sec. 2.1.2. To extract the rare Ξ_{cc}^{++} signal in pp collisions, the BDT is adopted, showing the high separation of signal and background performance.

The signal sample was generated by Delphes O2 Monte Carlo simulations [211], and the background sample used for combinatorial sets was generated by Pythia 8 in inelastic pp collisions. In the signal sample, 5000 Ξ_{cc}^{++} are injected per event in pp collisions at $\sqrt{s} = 14$ TeV, and the total number of events is 4000. The background sample is configured as an MB-triggered inelastic pp collision at $\sqrt{s} = 14$ TeV, and the total number of events is 2 Million.

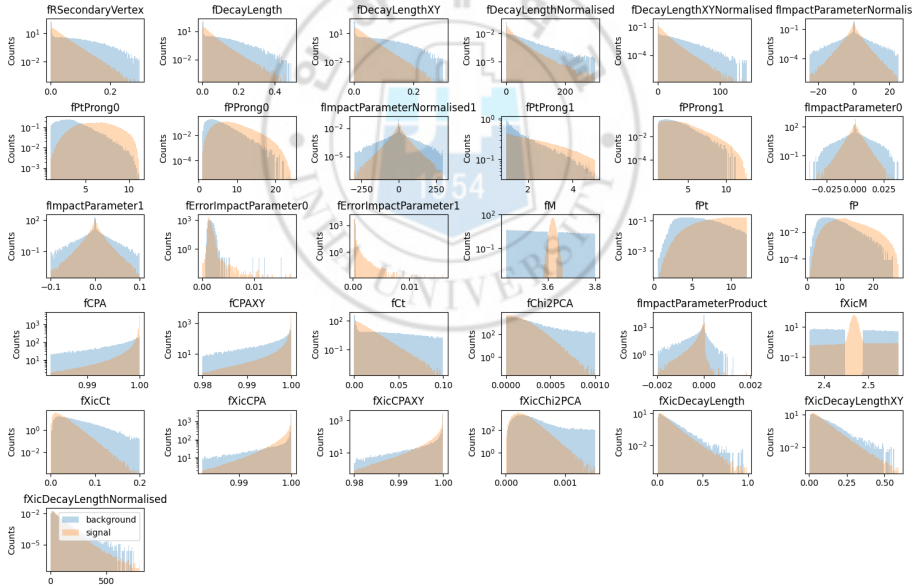


Figure 5.5: Signal and background distributions of each feature for $0 \leq p_T(\Xi_{cc}^{++}) < 12 \text{ GeV}/c$ ranges.

In the background sets, the Ξ_{cc}^{++} candidates are selected with the sideband regions of the invariant mass region. The mass of Ξ_{cc}^{++} is $3.621 \text{ GeV}/c$. Therefore, only the candidates with an invariant mass of $3.4 < m(\Xi_{cc}^{++}) < 3.601 \text{ GeV}/c$, or

Features	Cuts
$ \eta $	< 1.44
Radius from the secondary vertex	< 0.3
Decay length of Ξ_{cc}^{++}	< 0.5
Decay length of Ξ_{cc}^{++} in XY plane	< 0.3
Normalized decay length of Ξ_{cc}^{++}	< 300
Normalized decay length of Ξ_{cc}^{++} in XY plane	< 150
Impact parameter of Ξ_c^+	< 0.04
Impact parameter of $\pi^+(\leftarrow \Xi_{cc}^{++})$	< 0.1
Product of impact parameter	< 0.002
Normalized impact parameter of Ξ_c^+	< 25
Normalized impact parameter of $\pi^+(\leftarrow \Xi_{cc}^{++})$	< 300
Error of impact parameter	< 0.02
$c\tau$ of Ξ_{cc}^{++}	< 0.1
χ^2 of Ξ_{cc}^{++} PCA	< 0.001
Cosine of pointing angle to PV of Ξ_{cc}^{++}	> 0.983
Cosine of pointing angle to PV of Ξ_{cc}^{++} in XY plane	> 0.98
Cosine of pointing angle to PV of Ξ_c^+	> 0.983
Cosine of pointing angle to PV of Ξ_c^+ in XY plane	> 0.98
Invariant mass of Ξ_{cc}^{++}	> 3.4 and < 3.8
Decay length of Ξ_c^+	< 1
Decay length of Ξ_c^+ in XY plane	< 0.6
Normalized decay length of Ξ_c^+	< 800
$c\tau$ of Ξ_c^+	< 0.1
χ^2 of Ξ_c^+ PCA	< 0.0015
Invariant mass of Ξ_c^+	> 2.367 and < 2.567
$ n\sigma_{\text{TOF}} $	< 3

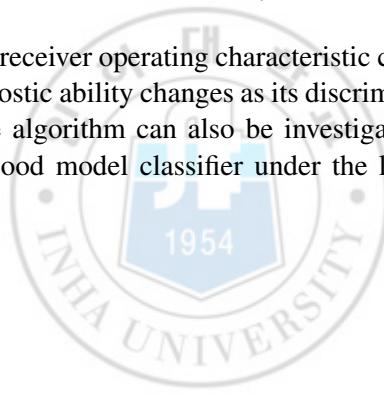
Table 5.1: pre-selections for Ξ_{cc}^{++} reconstruction.

$3.641 < m(\Xi_{cc}^{++}) < 3.8 \text{ GeV}/c$ are considered. The background sets are built by combining daughter particle candidates from the background sample. The pre-selections related to the decay topology are applied to reduce the sample size and computing time. The pre-selections are listed in Table 5.1. The pre-selections are defined based on comparing the distributions for signal and background. Each feature's signal and background distributions are shown in Fig. 5.5. In Fig. 5.5, a discrepancy between signal and background distributions on some features appears. This discrepancy is

ples. In the cross-validation step, the original sample is separated into n parts called folds. For each set of hyperparameters, $n-1$ folds are used for the optimization in training, and the remaining one is used for the test. This procedure is repeated after permuting the folds, and the mean value of all the permutations becomes a final result. In `hipe4ml`, the hyperparameters are automatically updated after their optimization. In this analysis, the number of DTs is assigned as from 200 to 1000 with a maximum of 5 folds, and the learning rate is assigned from 0.01 to 0.1. The training and testing are done for each p_T bin of Ξ_{cc}^{++} .

The training results can be obtained by distributions of the BDT score for the training and the test sets for each p_T bin in Fig. 5.7. The distributions of the model scores obtained from the test set agree with the training set, indicating that model has been appropriately trained. The result of the model, the BDT score, is applied to the full data sets as a selection. The event is accepted if it is above a threshold value of 0.999 for the whole p_T interval in this analysis to maximize the significance of Ξ_{cc}^{++} .

The ROC curve, the receiver operating characteristic curve, shows how a binary classifier system's diagnostic ability changes as its discrimination threshold change. The performance of the algorithm can also be investigated with the ROC curve, which characterizes a good model classifier under the large ROC curve area, as shown in Fig. 5.8.



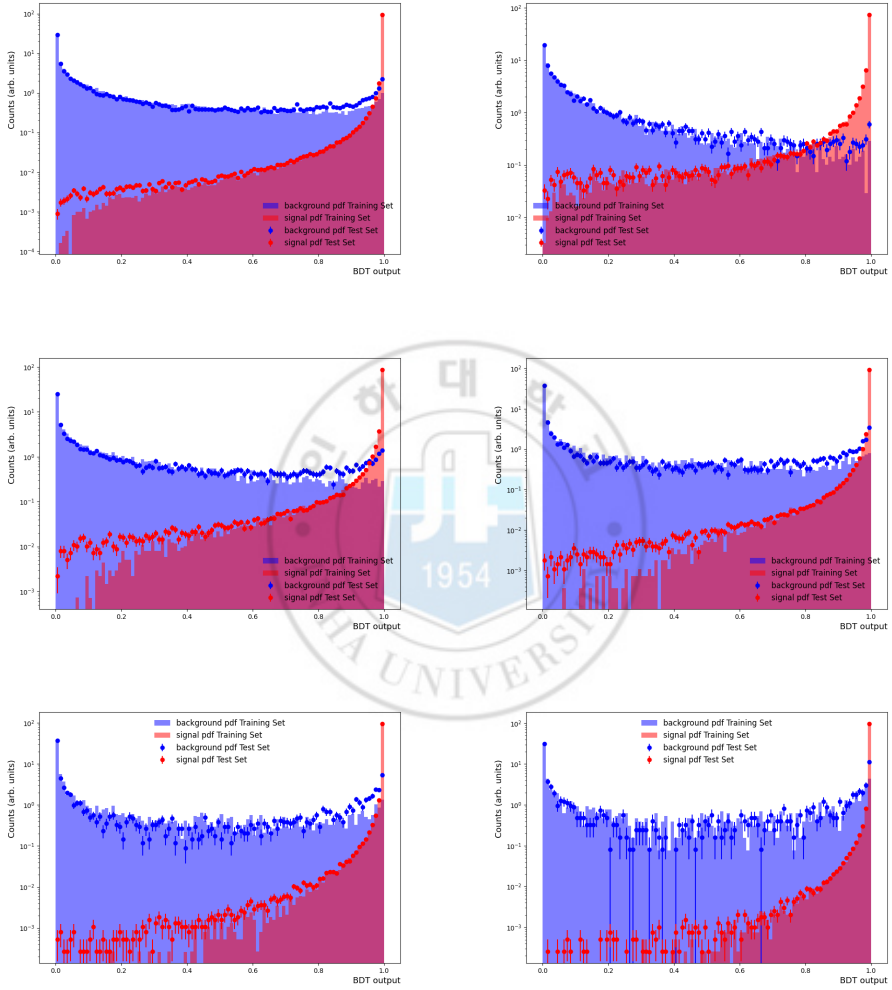


Figure 5.7: BDT output score distributions for signal and background used in the training and testing for each p_T bin.

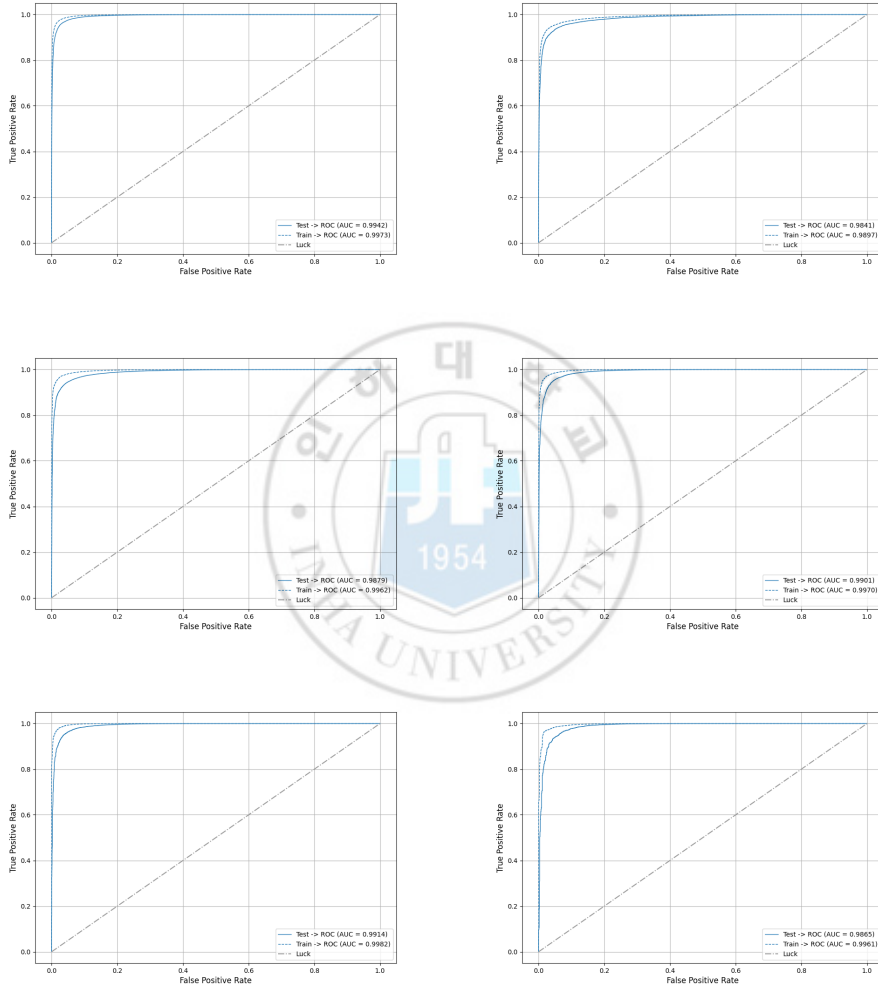


Figure 5.8: ROC curves from the training and test of the BDT models.

5.4 Results

The properties of the charm and multi-charm baryons considered in this study are reported in Table 5.2.

Particle	Mass (GeV/ c)	$c\tau$ (μm)	Decay channel	Branching Ratio (%)
Ξ_{cc}^{++}	3.621	76	$\Xi_c^+ + \pi^+$	5.0 (assumed)
Ξ_c^+	2.468	137	$p + K^- + \pi^+$	$(6.2 \pm 3.0) \times 10^{-3}$

Table 5.2: Particles and decay channels used in the reconstruction of the Ξ_{cc}^{++} . Where no measurement is available, a branching ratio of 5% is assumed. Numbers from [49].

Multi-charm baryon candidates are reconstructed by combining charm baryon candidates with pions, followed by the decay chain shown in Table 5.2. Selection criteria described in Table 5.1 and the threshold value of BDT output (0.999) are applied to the Ξ_{cc}^{++} candidates from both signal and background samples. In particular, pion, kaon, and proton identification are performed using the inner TOF detector for all decay products. Effectively, separating primary and secondary pions is allowed, as the latter arrives at the TOF detector with a delay. In this study, ML shows significant improvements in extracting the signal and rejecting the background with respect to the standard analysis.

5.4.1 Significance of $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ decay

The significance optimization is finding the best settlement between rejecting the background and keeping the signal. This study considers all prompt Ξ_{cc}^{++} baryons since they are directly injected in the dedicated MC production. The expected signal yield per event is calculated as follows:

$$S_{\Xi_{cc}^{++}}^{\text{pp}, 14\text{TeV}}(3\sigma, p_T) = \sigma_{\Xi_{cc}^{++}}^{\text{pp}, 14\text{TeV}} \cdot (acc \times \epsilon) \cdot 2\Delta p_T \Delta y \cdot \text{BR} \quad (5.3)$$

where the $\sigma_{\Xi_{cc}^{++}}^{\text{pp}, 14\text{TeV}}$ is expected cross section of Ξ_{cc}^{++} in pp collisions, based on the Ref. [214]. Based on Pythia calculations of charm baryons, it is assumed that 33% of the total yield is seen in the rapidity range $|y| < 1.44$, so additional scaling is applied to $\sigma_{\Xi_{cc}^{++}}^{\text{pp}, 14\text{TeV}}$. The $(acc \times \epsilon)$ is the expected efficiency of $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ obtained from the dedicated MC simulation, shown in Fig. 5.9. The bin width of p_T and y are considered, and the branching ratio of $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ is multiplied, based on Table 5.2.

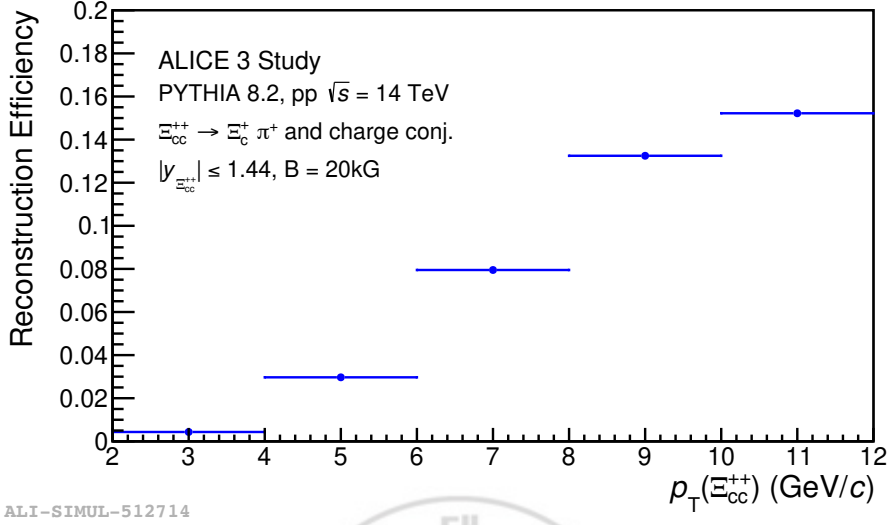


Figure 5.9: Expected reconstructed efficiency of $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ in pp collisions at $\sqrt{s} = 14$ TeV.

The background yield per event within 3σ ($B_{\Xi_{cc}^{++}}^{pp, 14\text{TeV}}(3\sigma, p_T)$) with respect to the signal mean is estimated from the MB-triggered inelastic pp collision events. The background sample does not include the injected Ξ_{cc}^{++} decays. The number of background candidates is counted from the invariant mass distribution, which is fitted with a polynomial function. In detail, the 3σ is extracted from the Ξ_{cc}^{++} injected signal sample, and the background yield is counted in the mass range of $3.591 < m(\Xi_c^+ + \pi^+) < 3.651 \text{ GeV}/c^2$. And the number of background candidates is normalized with the number of pp collisions (2 Million) to get a background yield per event.

The significance is calculated considering Ξ_{cc}^{++} baryons in $2 < p_T < 12 \text{ GeV}/c$ p_T range assuming an integrated luminosity of pp collisions of $\mathcal{L}_{\text{int}} = 18 \text{ fb}^{-1}$ at $\sqrt{s} = 14$ TeV with the inelastic cross section of $\sigma^{pp, 14\text{TeV}} = 70 \text{ mb}$, as Eq. [5.4](#).

$$\text{Significance}(3\sigma) = \frac{S_{\Xi_{cc}^{++}}^{pp, 14\text{TeV}}(3\sigma, p_T)}{\sqrt{S_{\Xi_{cc}^{++}}^{pp, 14\text{TeV}}(3\sigma, p_T) + B_{\Xi_{cc}^{++}}^{pp, 14\text{TeV}}(3\sigma, p_T)}} \cdot \sqrt{\sigma^{pp, 14\text{TeV}} \cdot \mathcal{L}_{\text{int}}} \quad (5.4)$$

where the $\sigma^{pp, 14\text{TeV}} \cdot \mathcal{L}_{\text{int}}$ corresponds to the number of events in pp collisions at $\sqrt{s} = 14$ TeV.

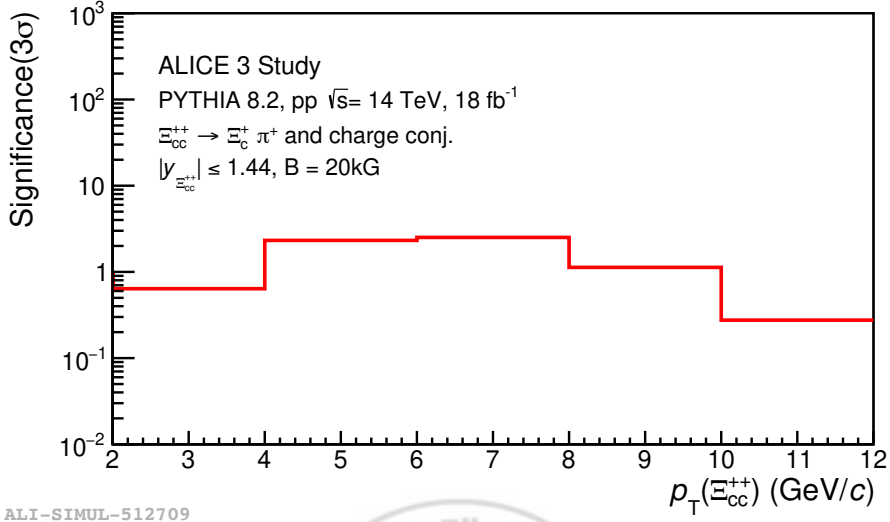


Figure 5.10: Expected significance of $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ in pp collisions at $\sqrt{s} = 14$ TeV.

The evaluated performance study using the decay channel $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ proceeds, the significance of the $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ in pp collisions are shown in Fig. 5.10. The significance shown in Fig. 5.10 is obtained in the worst condition, which means consider the lowset the branching ratio of the $\Xi_c^+ \rightarrow p + K^- + \pi^+$ channel. The branching ratio of $\Xi_c^+ \rightarrow p + K^- + \pi^+$ channel is argued as $0.45 \pm 0.21 \pm 0.07\%$ in Ref. [215], $1.135 \pm 0.002 \pm 0.387\%$ in Ref. [216], and as $2.2 \pm 0.8\%$ in Ref. [217]. Despite that, under the worst scenario, the significance of around 3 for $p_T > 4 \text{ GeV}/c$ is currently projected with $\mathcal{L}_{\text{int}} = 18 \text{ fb}^{-1}$ of pp collisions.

Chapter 6

Model calculation of bottomonium production

Quarkonium production has been studied extensively in a relativistic heavy-ion collision to understand the characteristic of the quark-gluon plasma, thanks to its large mass of the constituent quarks. The response to the medium produced by the collisions affects the physics and yield of the quarkonium state. Several models that consider the effects of dissociation and regeneration on the quarkonium state can explain the experimental results on the yield modification in heavy-ion collisions relative to pp collisions. Small collision systems like p—Au and p—Pb collisions have also shown yield modification. It can be explained by the beginning and end state effects or by the Cold Nuclear Matter effect (CNM), according to Ref. [218]. A model study in different small collision systems, such as pp, p-Pb, p-O, and O-O collisions, will aid in the quantitative understanding of nuclear effects on the production of $\Upsilon(nS)$ and be used to explore the yield modification on $\Upsilon(nS)$ in a small collision system. The suppression of $\Upsilon(1S)$ in Pb—Pb collisions is well described by a theoretical calculation evaluating the gluo-dissociation, inelastic parton scattering, and their inverse response [219].

A Monte-Carlo simulation framework was designed to integrate the medium created in heavy-ion collisions more realistically with event-by-event initial collision geometry and hydrodynamic evolution, which is introduced in Sec. 6.1. The Monte-Carlo framework is extended to small collision systems to study the medium response on the $\Upsilon(nS)$ states. The nuclear modification factor of $\Upsilon(nS)$ as a function of charged particle multiplicity ($dN_{\text{ch}}/d\eta$) and transverse momentum, and also the elliptic flow of $\Upsilon(nS)$ will be shown in Sec. 6.2.

6.1 Simulation for Heavy IoN Collision with Heavy-quark and ONia (SHINCHON)

This section introduces the SHINCHON simulation framework, which is specifically designed to simulate the medium response of quarkonium in the heavy-ion collision and small collision systems. The three components that construct the SHINCHON framework include the initial collision systems, which are regarded as the background medium and is produced by the MC-Glauber and hydrodynamics simulation. And when the dissociation effect is taken into account, the second component is regarded as the medium response of quarkonia. The modification of Upsilon is derived in the final section while taking feed-down contributions into the account.

6.1.1 Initial collision system

The following method is taken into consideration to construct the collision system's initial conditions; it is explained in Ref. [220]. The MC-Glauber framework is used to obtain the initial energy distribution in the $x - y$ plane of a collision system, see the Ref. [7]. The MC-Glauber setup is as follows: a nucleon-nucleon inelastic cross section is employed for small collision systems as 72 mb [221] at $\sqrt{s_{NN}} = 8$ TeV, the LHC energies. Additionally, the energy deposition of each nucleon involved in at least one inelastic collision is calculated using a gaussian with a width of $\sigma = 0.4$ fm. The MC-Glauber description's finer aspects were introduced in Sec 1.3.1. For the hydrodynamic simulation and a hadronic cascade simulation (SONIC), the deposited energy distribution from all damaged nucleons in each event is transformed into energy density. The scale factor is defined to match the measurement of the charged particle multiplicity ($dN_{ch}/d\eta$) at mid-rapidity in p-Pb collisions at $\sqrt{s_{NN}} = 8.16$ TeV [222, 223]. The scale factor is used for all events in a given collision system, by assuming that the scale factor does not significantly change in the collision systems with a similar number of participants. Therefore, the same scale factor is employed for p-Pb, p-O, and O-O collisions at $\sqrt{s_{NN}} = 8$ TeV. After the SONIC hydrodynamic simulation, the temperature information from the collision time to the freeze-out time for each step is stored based on the energy density of the medium. And the hadronic cascade simulation generates the properties of particles based on the final medium information after the freeze-out temperature.

The multiplicity distributions of charged particles in unbiased (0–100%) events in p-O, O-O, and p-Pb collisions at $\sqrt{s_{NN}} = 8$ TeV are shown in the left panel of Fig. 6.1. To get multiplicity distributions for all small collision systems, the scale factor is adjusted to correspond to the measurement of multiplicity distribution in p-Pb collisions at $\sqrt{s_{NN}} = 8.16$ TeV. Additionally, the bulk viscosity ratio is set to

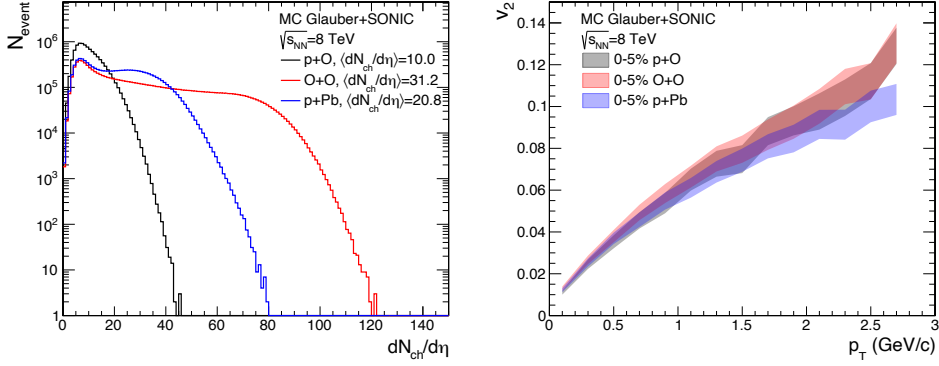


Figure 6.1: Left panel: charged particle multiplicity distribution in p-O (black), O-O (red), and p-Pb (blue) collisions at $\sqrt{s_{NN}} = 8$ TeV from MC-Glauber and SONIC. Right panel: the elliptic flow of the charged particles as a function of p_T in 0–5% central p-O (black), O-O (red), and p-Pb (blue) collisions at $\sqrt{s_{NN}} = 8$ TeV from SONIC. The uncertainties of the calculations indicate the statistical uncertainty of the corresponding simulated events.

be zero, and the viscosity to entropy density ratio is set to be 0.08 in the SONIC calculations. The mean charged particle multiplicity is also shown in the left panel of Fig. 6.1, for p-O is 10.0, for O-O is 31.2, and for p-Pb is 20.8. As seen in the right panel of Fig. 6.1, the elliptic flow of charged particles in 0–5% high multiplicity events is checked. In comparison to p-O and O-O collisions at higher p_T , the elliptic flow of charged particles in p-Pb collisions is slightly lower. Although there are different frameworks for initial conditions [224, 225], MC-Glauber is used in the SHINCHON framework in this study as an initial condition for a straightforward case.

6.1.2 Medium response on Upsilon

Temperature distributions for the medium response of $\Upsilon(nS)$ are maintained for each step of the SONIC simulation from the collision time to the freeze-out time. The x and y positions of the $\Upsilon(nS)$ are established based on the initial energy density distribution at the start of the second part, where $\Upsilon(nS)$ are formed in the medium. The measurement of the p_T distribution in Pb-Pb collisions at $\sqrt{s} = 5.02$ [19] yields the p_T distribution of $\Upsilon(.S)$. The Tsallis function [226] is used to fit the p_T distribution of $\Upsilon(1S)$ in $\sqrt{s} = 5.02$ [19], and p_x and p_y are determined based on a randomly assigned azimuthal angle. The p_T distribution of $\Upsilon(2S)$ and $\Upsilon(3S)$ is assumed to be

the same as that of $\Upsilon(1S)$. The initial p_T distribution of $\Upsilon(nS)$ is assumed to be the same for 5.02 TeV and 8 TeV. Also, possible difference in the p_T shape of $\Upsilon(nS)$ is expected to be negligible because the results are obtained in fine p_T bins.

The medium response is adopted based on the procedure described in Ref. [219]. In a partonic picture, quarkonium dissociation occurs through two scattering processes, gluo-dissociation ($\Upsilon + g \rightarrow b + \bar{b}$) and inelastic parton scattering ($\Upsilon + p \rightarrow b + \bar{b} + p$ with $p = g, q, \bar{q}$). The inverse process of gluo-dissociation and inelastic parton scattering is relevant to quarkonium regeneration. Dissociation, regeneration, and elastic scatterings describe the dynamic evolution of a quarkonium state. If the inverse distance between a heavy quark and its antiquark is larger than the temperature scale, the order of contribution of elastic collisions is higher than dissociation and regeneration [227]. Also, the regeneration effect in the small collision systems on the Upsilon state is expected to be negligible. Thus, neglecting these terms, the Boltzmann equation for quarkonium is given by [111].

$$\left(\frac{\partial}{\partial t} + v \cdot \frac{\partial}{\partial x} \right) f_{\Upsilon}(t, x, q) = -\Gamma_{\text{diss}}^{\text{gluo} + \text{inel}}(t, x, q) f_{\Upsilon}(t, x, q) \quad (6.1)$$

where $v = q/q^0$ is the velocity of the Upsilon states, and the distribution functions f_{Υ} contain the fugacity factors of γ_b^2 . The first term on the right-hand side represents the dissociation with the thermal width Γ_{diss} .

Using the thermal width, the fraction of survived Upsilon for a given time step (Δt) is calculated as:

$$\frac{N(t + \Delta t, p_T)}{N(t, p_T)} = e^{-\int_t^{t+\Delta t} dt' \Gamma_{\text{diss}}(t', p_T)} \quad (6.2)$$

where Γ_{diss} is the thermal width depending on the medium temperature and p_T of the Upsilon, and Δt , which is the time interval is set to 0.02 fm/c in the hydrodynamic simulation.

The model calculation obtained the thermal width from numerical calculations considering gluo-dissociation and inelastic parton scattering from the model calculation [219]. The medium temperature is from the SONIC simulation according to x , y , and t values. The dissociation temperature of $\Upsilon(nS)$ states is defined as 600 MeV, 240 MeV, and 190 MeV for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$, respectively [228, 229], which is the in-medium binding energy of quarkonium states reaches zero.

The left column of Fig. 6.2 shows the thermal width as a function of temperature for $\Upsilon(1S)$ (top), $\Upsilon(2S)$ (middle), and $\Upsilon(3S)$ (bottom). Each colored lines indicate a distribution for the given different p_T of the $\Upsilon(nS)$, and the black lines indicate for intermediate p_T of the $\Upsilon(nS)$. The thermal width generally increases with p_T and shows a larger value for excited states. The thermal width for $\Upsilon(3S)$ is calculated

in the limited p_T range by the model calculations, so the thermal width is estimated using the thermal width of $\Upsilon(2S)$ for $p_T > 2 \text{ GeV}/c$ as following:

$$\Gamma_{\text{diss}}^{\Upsilon(3S)}(p_T) = \Gamma_{\text{diss}}^{\Upsilon(3S)}(2 \text{ GeV}/c) \frac{\Gamma_{\text{diss}}^{\Upsilon(2S)}(p_T)}{\Gamma_{\text{diss}}^{\Upsilon(2S)}(2 \text{ GeV}/c)} \quad (6.3)$$

A formation time of $\Upsilon(nS)$ for each states was assigned as 0.5, 1.0, and 1.5 for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$, respectively, based on values in Ref. [230]. The dissociation at the pre-resonance stage is not considered, so the medium response on the Υ states is turned off for $\tau < \gamma\tau_{\text{form}}$, where γ is the Lorentz factor depends on the velocity of the $\Upsilon(nS)$ states. This configuration of the medium response can well reproduce the measurement in heavy-ion collisions and is used for small systems consistently. After the formation time, a survival rate is calculated in each time step for each Υ based on Eq. 6.2. The Υ survived when a random number from the uniform distribution within 0–1 is smaller than the survival rate from the Eq. 6.2. If the Υ survives, the position of the Υ for the next time step is calculated based on the momentum and time step (Δt). The calculation procedure is repeated for a decreased temperature at the new position until the temperature reaches a critical temperature of 170 MeV.

The number of 1000 events is generated by MC-Glauber and SONIC for the simulation result. The number of generated $\Upsilon(nS)$ for a given event is scaled with the number of inelastic collisions from the MC-Glauber. As mentioned before, the contribution of the regeneration effect is not considered, which is not predicted to be significant in small collision systems.

As a result of the medium response, the nuclear modification factor and elliptic flow are obtained from the simulation framework. The nuclear modification factor is used to investigate the yield modification, and the elliptic flow is used to investigate the medium geometry effect on Υ states.

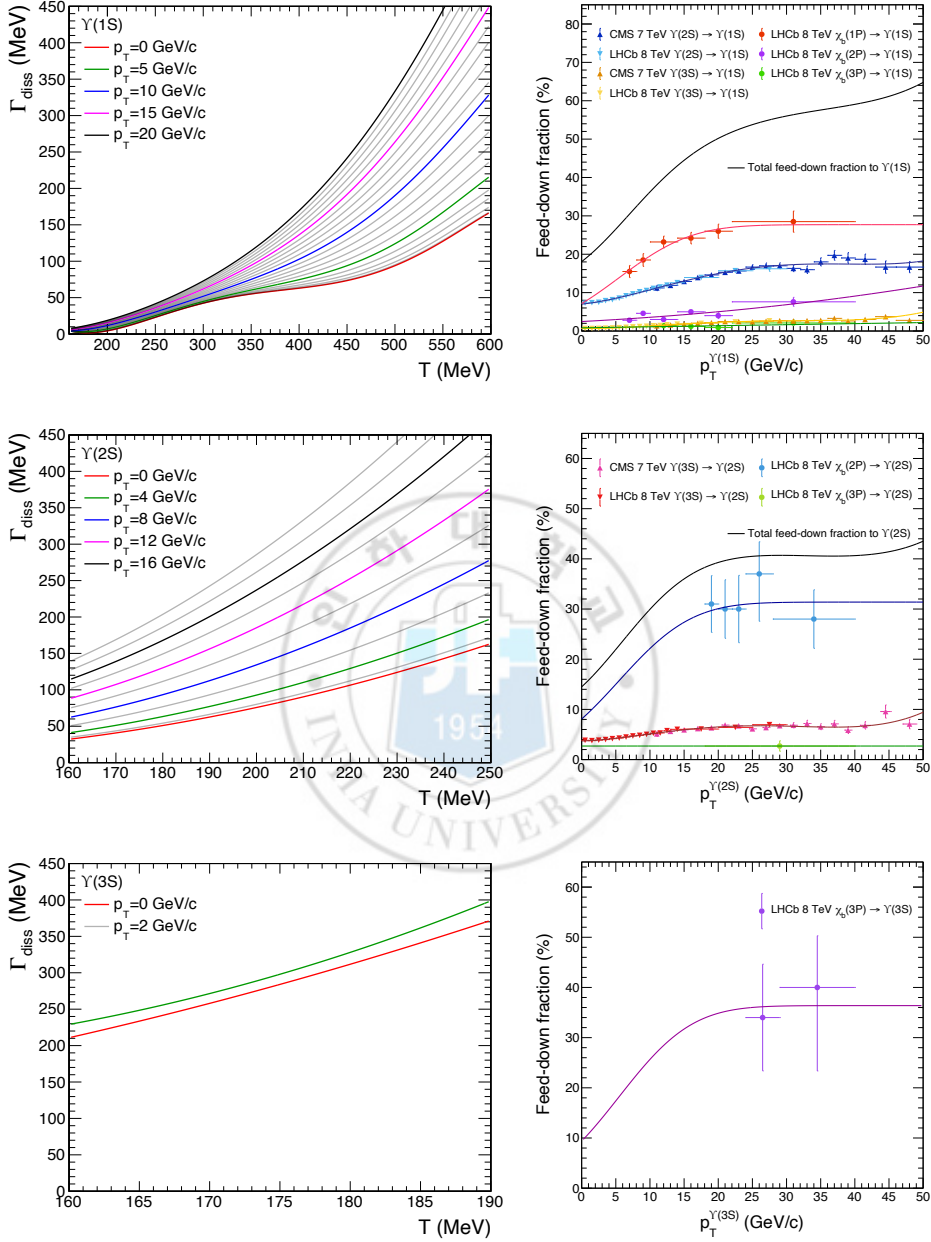


Figure 6.2: Left column: thermal width as a function of temperature for different p_T of $\Upsilon(1S)$ (top), $\Upsilon(2S)$ (middle), and $\Upsilon(3S)$ (bottom). Right column: feed-down fraction for $\Upsilon(1S)$ (top), $\Upsilon(2S)$ (middle), and $\Upsilon(3S)$ (bottom) states.

Nuclear modification factor

The nuclear modification factors (R_{pA} or R_{AA}) are defined as the ratios of the production cross sections of Υ in p-A or A-A collisions to the expected production cross sections extrapolated from pp collisions (see Sec. 1.4.7). To make an analogous comparison, the ratio of the number of the survived Υ s and the number of generated Υ s in the whole medium response simulation is treated as the nuclear modification factor.

The nuclear modification factor is calculated as a function of the multiplicity of charged particles and as a function of the p_T of the Υ .

Elliptic flow

The elliptic flow (v_2) is calculated from the azimuthal angle distribution of survived $\Upsilon(nS)$ with respect to the event plane angle (see Sec. 1.4.6).

$$\frac{dN}{d(\phi - \Psi)} \propto 1 + 2v_2(p_T) \cos(2(\phi - \Psi)), \quad (6.4)$$

where ϕ is the azimuthal angle of the $\Upsilon(nS)$, and Ψ is the event plane angle. The event plane angle for a given event is computed with the initial energy density information from the MC-Glauber.

The elliptic flow is calculated as a function of the p_T of the Υ in unbiased (0–100%) events and in the high-multiplicity (0–5%) events.

6.1.3 Feed-down contribution

The result from the simulation is the result of the prompt Υ . The contributions from feed-down decay, or decay from higher excited states, must be carefully considered in order to analyze the medium response of inclusive bottomonium production, which is measured by experimental groups.

The feed-down decay from the strong or electromagnetic modes of excited states is considered to evaluate the inclusive Υ production. On the inclusive yields, it is considered that the feed-down decay from the H , Z , and W bosons to Υ states has no impact. The feed-down contribution from a Q_m state to a Q_n state (Q_m being a higher excited state) can be calculated as:

$$\mathcal{F}_{Q_n}^{Q_m} = \mathcal{B}(Q_m \rightarrow Q_n) \frac{\sigma_{Q_m}}{\sigma_{Q_n}} \quad (6.5)$$

where $\mathcal{B}$ indicates the branching ratio of Q_m into Q_n and σ_Q is the production cross section of the given state.

To estimate the feed-down contribution, the published experimental measurements in pp collisions at LHC energies by CMS [231] and LHCb [232, 233] are considered. Fig. 6.2 shows the feed-down fraction from higher excited states for each $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ states. The feed-down fractions from CMS measurements are estimated with the cross section ratios multiplied by the branching ratios, published in PDG world-average [49] as in Eq. 6.5. The points from the measurements of each feed-down contribution are fitted with an empirical function. The solid black line shows the sum of the feed-down contribution, which is the total amount of feed-down fraction for the given Υ state. Because of the absence of experimental measurements at low- p_T for $\Upsilon(2S)$ and $\Upsilon(3S)$, the fit function is assumed to be the same as for the P-wave state feed-down function of $\Upsilon(1S)$. And then, the fit function is normalized to match the data points at high- p_T for each $\Upsilon(2S)$ and $\Upsilon(3S)$ states, respectively.

The feed-down contribution for the nuclear modification factor and elliptic flow of inclusive $\Upsilon(nS)$ is evaluated by computing a weighted average of the quantities for possible states as:

$$R_n(p_T) = \sum_i R_i(p_T) \mathcal{F}_{Q_n}^{Q_i}(p_T), \quad (6.6)$$

where R_n is the weighted averaged value such as modification factor and azimuthal distribution of a given $\Upsilon(nS)$ state. R_i is the value for a given state contributing the $\Upsilon(nS)$ state, and $\mathcal{F}_{Q_n}^{Q_i}$ is the feed-down fraction, which is shown in the right column of Fig. 6.2.

Given that the mother and daughter particles have similar masses and that the daughter particle mostly carries the momentum of the mother, the possibility of p_T changes by feed-down decays is not taken into account. The measurements of the χ_b states are not enough to consider the feed-down decays on $\Upsilon(nS)$ states. Therefore, it is assumed that $R_{\Upsilon(2S)} \approx R_{\chi_b(1P)}$ and $R_{\Upsilon(3S)} \approx R_{\chi_b(2P)} \approx R_{\chi_b(3P)}$ like the study in Ref. [219].

6.2 Results

This section presents the obtained results of nuclear modification factors (R_{pA} , R_{AA}) and elliptic flow (v_2) for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ mesons in p-Pb, p-O, and O-O collisions at $\sqrt{s_{NN}} = 8$ TeV. The contribution from the feed-down decays is corrected in all calculations as described in Sec. 6.1.3.

The performance of our framework is demonstrated in Pb-Pb collisions and compared with the experimental measurements at LHC as described in Sec. 6.2.1.

The results of the nuclear modification factors and v_2 in the three small collision systems are described in Sec. 6.1.2. Also, The comparison of the experimental measurements in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV and the result of model calculation in a small collision system for nuclear modification factor is described.

6.2.1 Framework demonstration in Pb-Pb collisions

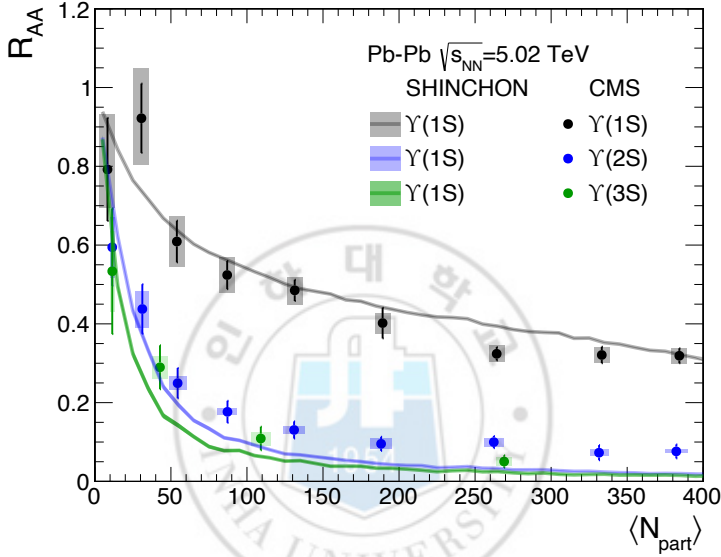


Figure 6.3: Calculated R_{AA} for Υ states as a function of $\langle N_{part} \rangle$ in Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV from the simulation. The uncertainties of the calculations indicate the statistical uncertainty of the corresponding simulated events. The data points are taken from [19, 42].

The medium response on Υ in full simulation is demonstrated in Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. Fig. 6.3 shows the R_{AA} distributions for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ as a function of $\langle N_{part} \rangle$ in Pb-Pb collision at $\sqrt{s_{NN}} = 5.02$ TeV. The simulation result is compared with the measurements from CMS [19, 42] in Pb-Pb collision at $\sqrt{s_{NN}} = 5.02$ TeV. The calculation from the simulation well describes the experimental measurement for $\Upsilon(1S)$ after applying feed-down corrections in all $\langle N_{part} \rangle$ region and for $\Upsilon(2S)$ and $\Upsilon(3S)$ in small $\langle N_{part} \rangle$ region. However, the calculations from the simulation underestimate the experimental measurements for $\Upsilon(2S)$ and $\Upsilon(3S)$ in large $\langle N_{part} \rangle$ region. Therefore, the excited Υ calculations show consistency with the measurements in peripheral collisions, although discrepancy is

found towards central Pb-Pb collisions. The discrepancy in central collisions might come from the exclusion of regeneration processes, which is expected to be more noticeable in the large $\langle N_{\text{part}} \rangle$ region.

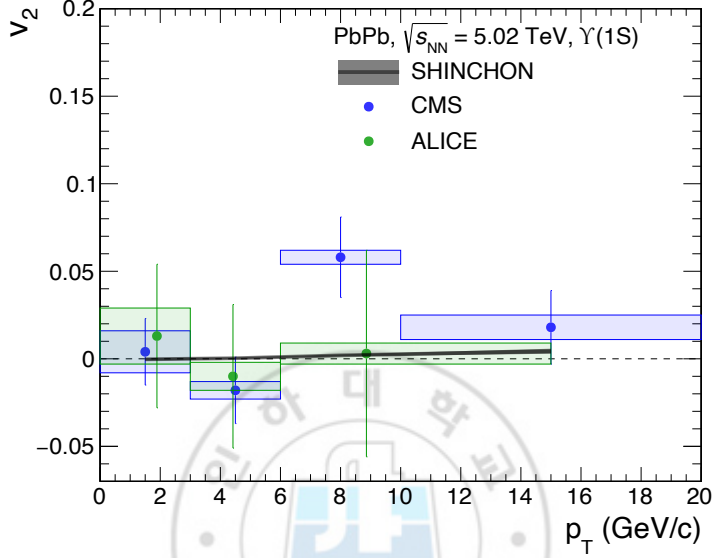


Figure 6.4: Calculated v_2 for $\Upsilon(1S)$ state as a function of p_T in Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The uncertainties of the curves represent the statistical uncertainty of the corresponding simulated events. The data points are taken from results by ALICE [43] and CMS [44].

Fig. 6.4 shows the calculations from the simulation of v_2 , as a function of p_T for $\Upsilon(1S)$ mesons. The v_2 result is obtained under the same conditions for the Pb-Pb R_{AA} mensuration. The v_2 result shows very small v_2 values in Pb-Pb collisions for $\Upsilon(1S)$ mesons, which is consistent with the measurements from ALICE [43] and CMS [44]. The small v_2 values might suggest that the $\Upsilon(1S)$ meson is not largely affected by the medium pressure due to its large mass.

The results from the SHINCHON framework well describe the experimental measurements in Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV even without the regeneration processes of medium response. In small collision systems, the regeneration processes can be negligible. Therefore, calculations from the SHINCHON framework in small collision systems can derive reasonable results.

6.2.2 Nuclear modification factor in small collision systems

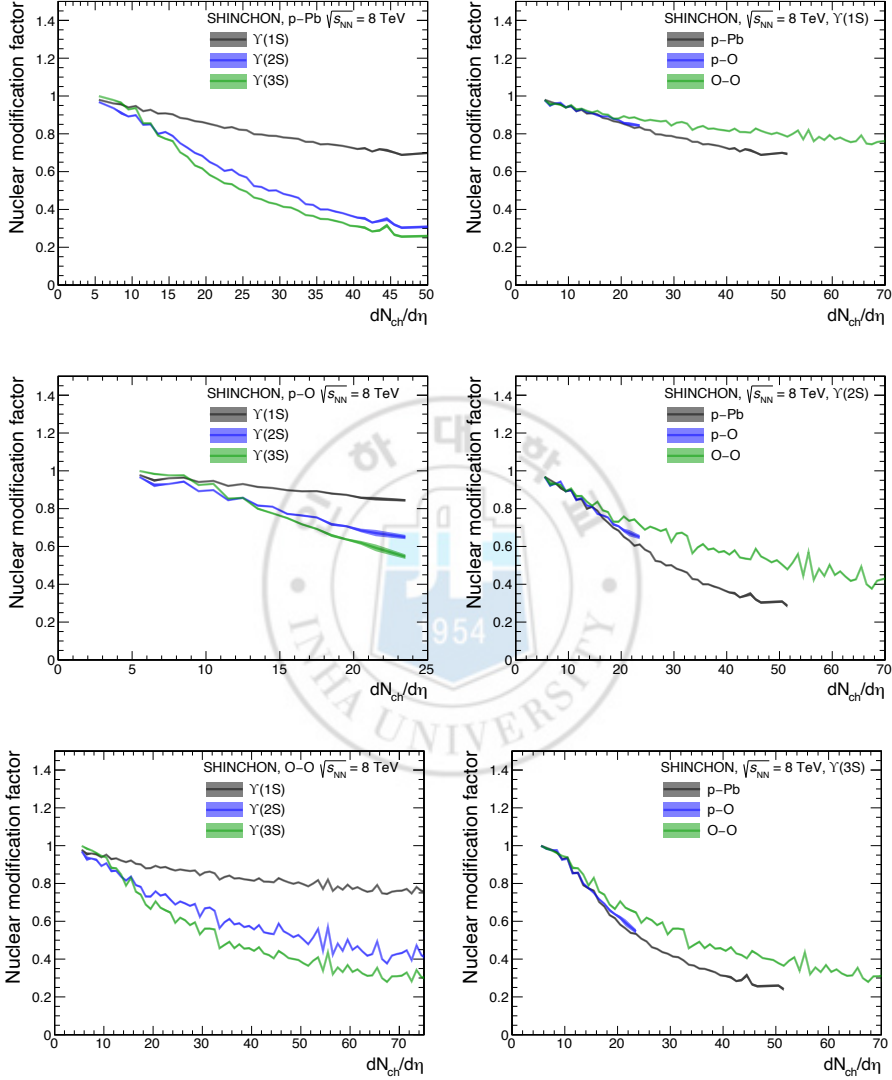


Figure 6.5: Left column: nuclear modification factor of $\Upsilon(1S)$, $\Upsilon(2S)$, $\Upsilon(3S)$ as a function of $dN_{ch}/d\eta$ in p-Pb (top), p-O (middle), and O-O (bottom) collisions at $\sqrt{s_{NN}} = 8$ TeV. Right column: nuclear modification factor of $\Upsilon(1S)$ (top), $\Upsilon(2S)$ (middle), $\Upsilon(3S)$ (bottom) as a function of $dN_{ch}/d\eta$ in p-Pb, p-O, and O-O collisions at $\sqrt{s_{NN}} = 8$ TeV. The uncertainties represent the statistical uncertainty.

The calculated nuclear modification factor as a function of $dN_{\text{ch}}/d\eta$ for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ mesons in p-Pb, p-O, and O-O collisions at $\sqrt{s_{\text{NN}}} = 8$ TeV are shown in the left column of Fig. 6.5. For each collision system, the black line represents the R_{pA} of $\Upsilon(1S)$, and the blue and green line represents the R_{pA} of $\Upsilon(2S)$ and $\Upsilon(3S)$, respectively. The R_{pA} of Υ gradually decreases with increasing event multiplicity for all three Υ states in all three small collision systems. This behavior suggests that there is more suppression in high multiplicity events than low multiplicity events due to the higher density of the medium. Also, the amount of suppression towards higher multiplicity events is sequentially ordered following the according to the related Υ binding energy ($E(\Upsilon(1S)) > E(\Upsilon(2S)) > E(\Upsilon(3S))$). The probability of interaction breaking in Υ states increases when the binding energy decreases. Therefore, there is more suppression in $\Upsilon(3S)$ states compare with $\Upsilon(1S)$ and $\Upsilon(2S)$.

To investigate the amount of modification among different small collision systems, the nuclear modification factors are separately shown for each Υ state as shown in the right column of Fig. 6.5. For each Υ states, the black line represents the R_{pA} in p-Pb collisions, and the blue and green line represents the R_{pA} in p-O collisions and in O-O collisions, respectively. In the low multiplicity region ($dN_{\text{ch}}/d\eta < 25$), the suppression of all Υ states is similar in p-Pb and p-O collisions, while a slightly less suppression is shown in O-O collisions in particular for $\Upsilon(2S)$ and $\Upsilon(3S)$. This trend stands out in the high multiplicity region. At the higher multiplicity region, where the results in p-Pb and O-O collisions are compared, Υ states are more robustly suppressed in p-Pb collisions than in O-O collisions which is consistent with the finding at low multiplicity. The system size of the medium is the largest in O-O collisions and smallest in p-O and p-Pb collisions. The SHINCHON simulation results show that the suppression amount is not proportional to the size of the created system. Even though, the system size is larger in a given multiplicity event in O-O collision. However, on the other hand, the energy density of the medium in O-O collision is lower than p-Pb collision. Therefore, it can be concluded that the weaker suppression observed in O-O collisions is due to the smaller energy density compared to p-O and p-Pb collisions.

Fig. 6.6 shows the same quantities as a function of p_T . In the left column of Fig. 6.6, the ordering of the nuclear modification factors for all collision systems is similar at low- p_T for all Υ states as in Fig. 6.5 while reversed R_{pA} between $\Upsilon(2S)$ and $\Upsilon(3S)$ states is found at high- p_T region. For each collision system, the black line represents the R_{pA} of $\Upsilon(1S)$, and the blue and green line represents the R_{pA} of $\Upsilon(2S)$ and $\Upsilon(3S)$, respectively. Similarly, In the right column of Fig. 6.6, the ordering of the nuclear modification factors for all Υ states is similar at low- p_T in all small collision systems as in Fig. 6.5 while reversed R_{pA} between p-Pb and O-O collisions is found at high- p_T region. For each Υ state, the black line represents

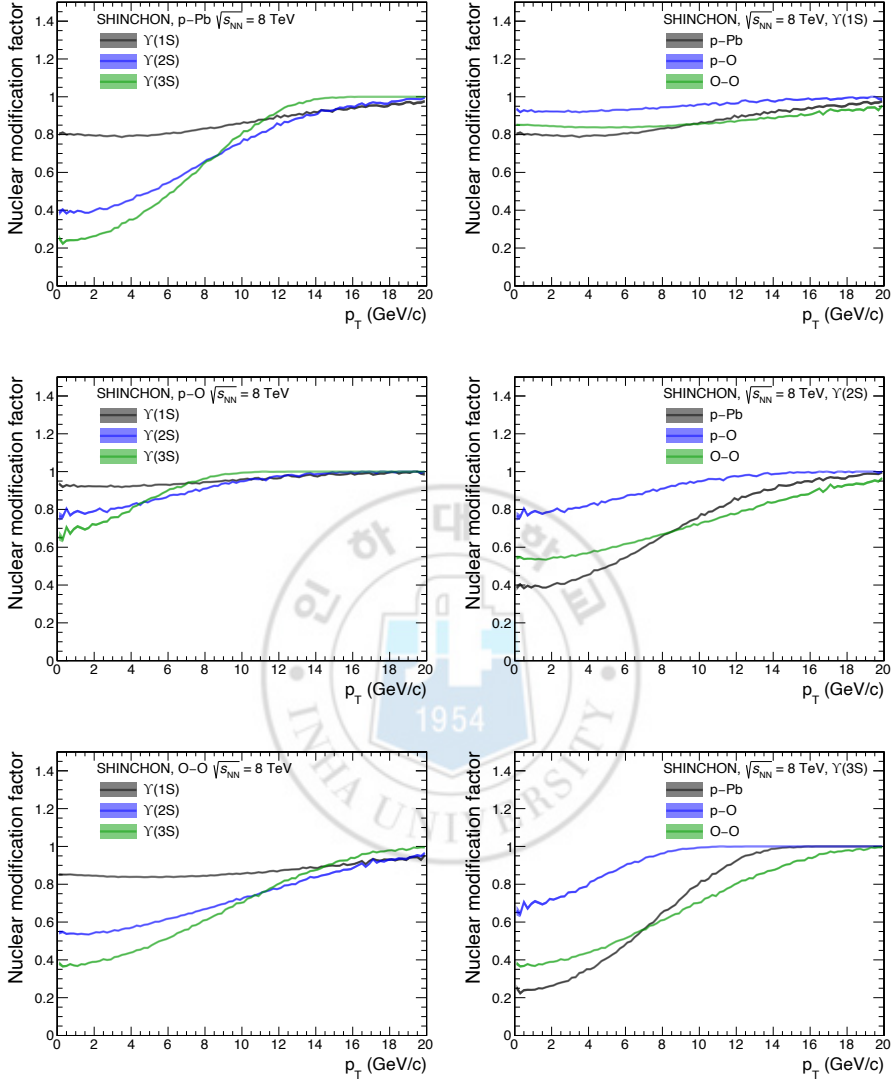


Figure 6.6: Left column: nuclear modification factor of $\Upsilon(1S)$, $\Upsilon(2S)$, $\Upsilon(3S)$ as a function of p_T in p-Pb (top), p-O (middle), and O-O (bottom) collisions at $\sqrt{s_{NN}} = 8$ TeV. Right column: nuclear modification factor of $\Upsilon(1S)$ (top), $\Upsilon(2S)$ (middle), $\Upsilon(3S)$ (bottom) as a function of p_T in p-Pb, p-O, and O-O collisions at $\sqrt{s_{NN}} = 8$ TeV. The uncertainties represent the statistical uncertainty.

the R_{pA} in p-Pb collisions, and the blue and green line represents the R_{pA} in p-O collisions and in O-O collisions, respectively. The formation time of the Υ states is delayed towards the higher p_T region due to relativistic motion. Therefore, each state's effective interaction time decreases between medium and Υ , reminding that the formation time of Υ states is determined as 0.5, 1.0, and 1.5 for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$, respectively. This behavior is more promising in p-Pb collisions because of the smaller initial medium size and the largest medium energy density.

The Υ suppression is affected by the medium size and energy density in the opposite direction for low- and high- p_T . These results suggest that the trend of the p_T -dependent nuclear modification factor is sensitive to the formation time and initial collision geometry.

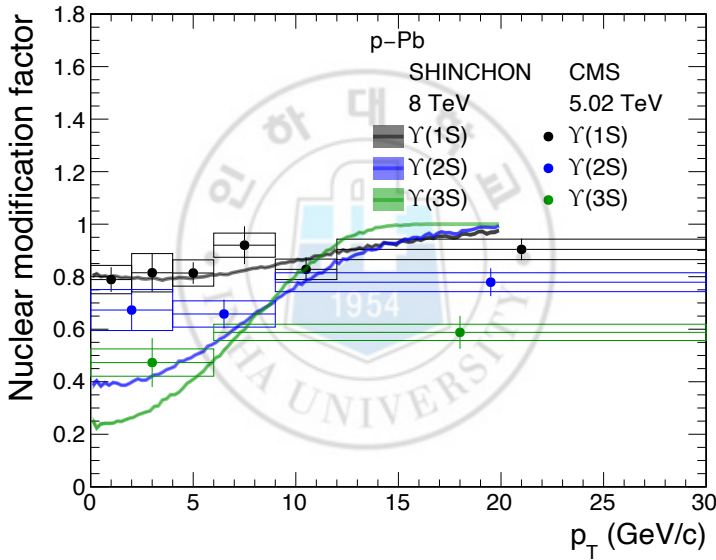


Figure 6.7: Calculated nuclear modification factor for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ as a function of p_T in p-Pb collisions at $\sqrt{s_{NN}} = 8$ TeV compared with the CMS results at $\sqrt{s_{NN}} = 5.02$ TeV.

The results from the SHINCHON simulation in p-Pb collisions at $\sqrt{s_{NN}} = 8$ TeV are compared with the measurements of R_{pA} as a function of p_T measured by CMS in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [234] as shown in Fig. 6.7. The results agree with the measurements for $\Upsilon(1S)$, while the discrepancies are observed for $\Upsilon(2S)$ and $\Upsilon(3S)$. The average multiplicity at $\sqrt{s_{NN}} = 8$ TeV is about 15% higher than that at 5.02 TeV [222]. At low- p_T , when Υ mainly experiences the entire medium evo-

lution, the multiplicity dependence of collision energy has a greater impact on the modification of Υ yields. Our calculations still overestimate the measurements even if the mean p_T of the highest p_T bin of the CMS data for the $\Upsilon(2S)$ and $\Upsilon(3S)$ meson is expected to be slightly lower than the center of the bin. These inconsistencies are possibly related to the formation time of Υ states and the initial collision geometry of the collision systems, as discussed in Fig. 6.6. According to model calculations for an initial-state model that takes energy loss and modified nuclear parton distribution functions into account, $\Upsilon(1S)$ is suppressed about 10% suppression of $\Upsilon(1S)$ at mid-rapidity [222, 235, 218].

6.2.3 Anisotropic flow in small collision systems

The v_2 is calculated in the same configuration as in the calculation for the nuclear modification factors in Sec. 6.1.2 and shown in Fig. 6.8. For all three Υ states, the v_2 values are consistent with zero in the overall p_T region in p-Pb, p-O, and O-O collisions. Also, their v_2 are very similar among the three small collision systems. It may suggest that the extended formation time of Υ s towards high- p_T is not as significant as implemented in the operated model.

To investigate the relation between v_2 and the amount of suppression, the v_2 values were computed using the azimuthal angle distribution of $\Upsilon(nS)$ in the 0–5% high-multiplicity events. For all Υ states, the v_2 values are more considerable at high multiplicity events only with a small magnitude.

Even while excited states and higher multiplicity events have more suppression for Υ meson yields, their v_2 values remain constant. Due to the low- p_T *Upsilon* mesons' slow traversal, they are unable to leave the medium before they reach the chemical freeze-out temperature. Despite their considerable suppression, such *Upsilon* mesons are unable to capture the anisotropy of the initial collision geometry. It is known that the path-length dependent suppression may result in a non-zero v_2 that can be investigated for Υ mesons at high p_T . However, with current simulation configurations, there is no significant non-zero v_2 that deviates from that at low- p_T in small collision systems for all Υ states.

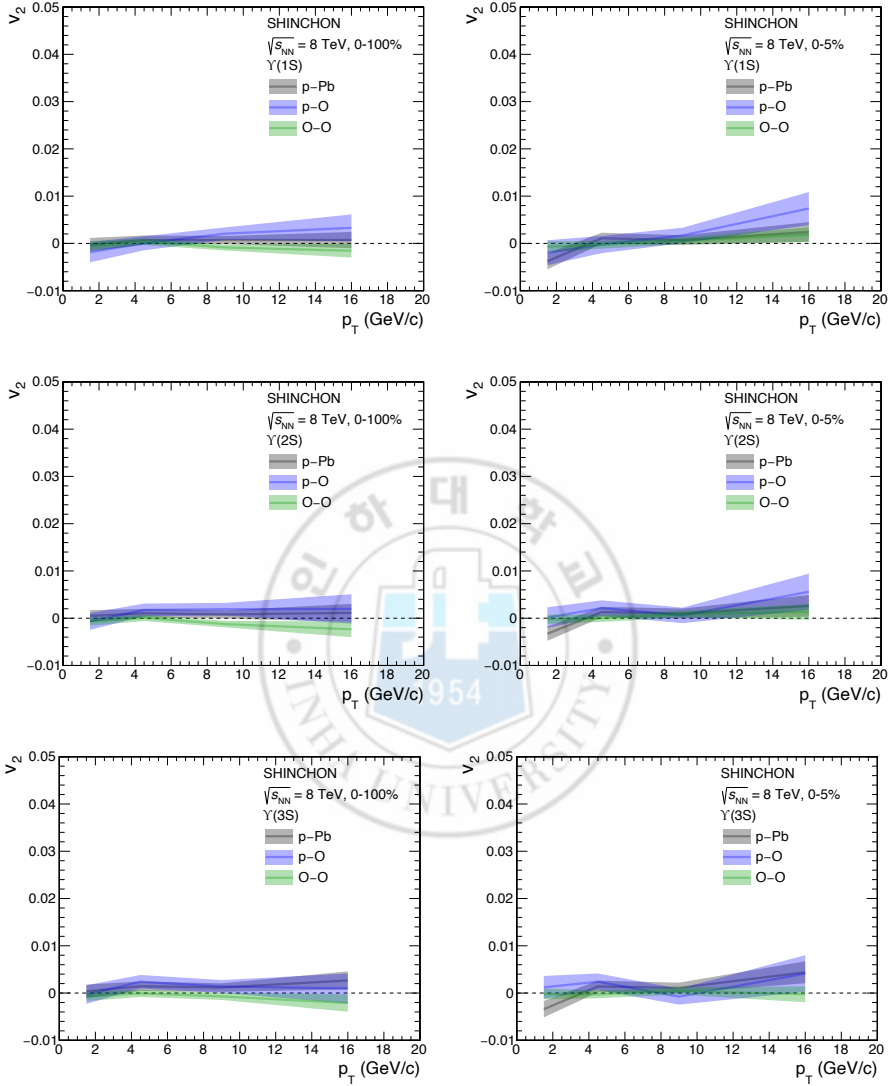


Figure 6.8: Left column: calculated v_2 in p-Pb, p-O, and O-O collisions as a function of p_T for $\Upsilon(1S)$ (top), $\Upsilon(2S)$ (middle), $\Upsilon(3S)$ (bottom) at $\sqrt{s_{NN}} = 8$ TeV. Right column: calculated v_2 in p-Pb, p-O, and O-O collisions as a function of p_T for $\Upsilon(1S)$ (top), $\Upsilon(2S)$ (middle), $\Upsilon(3S)$ (bottom) at $\sqrt{s_{NN}} = 8$ TeV in 0-5% high multiplicity events. The uncertainties come from the fitting procedure which is affected by the fluctuation of the angular distribution.

Chapter 7

Summary and outlook

The measurement of the prompt charm-strange baryons Ξ_c^0 at midrapidity in pp collisions at $\sqrt{s} = 13$ TeV is introduced in Chapter 4. The transverse momentum (p_T) dependence of the ratio of Ξ_c^0 and the D^0 , and ratio of the Ξ_c^0 and the $\Sigma_c^{0,+,++}$ and Λ_c^+ is discussed. The charming baryon-to-meson ratio increases towards low p_T up to around 0.3, and converges to 0.1, which is consistent with model calculations considering the fragmentation function measured in e^+e^- and ep collisions. The Ξ_c^0/Λ_c^+ ratio is approximately 0.5, and $\Xi_c^{0,+}/\Sigma_c^{0,+,++}$ ratio is approximately 2, showing no strong p_T dependence within the current uncertainties. Model calculation considering the coalescence mechanism near the $p_T = 0$ describes better the charmed baryon-to-hadron ratio yield in measured p_T interval. The branching fraction ratio $\text{BR}(\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e)/\text{BR}(\Xi_c^0 \rightarrow \Xi^- \pi^+) = 1.38 \pm 0.14 (\text{stat}) \pm 0.22 (\text{syst})$ is obtained with a total uncertainty reduced by a factor 3 concerning the global average reported by the PDG [49]. The measurement of Ξ_c^0 can provide strong constraints to the model calculations, comparing with the model calculations, and also provide evidence that different charm hadronization processes e^+e^- and pp collisions.

The evaluated performance study using the decay channel $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ in Run 5 with ALICE 3 proceeds in Chapter 5. The significance of the $\Xi_{cc}^{++} \rightarrow \Xi_c^+ \pi^+ \rightarrow (pK^- \pi^+) \pi^+$ in pp collisions is discussed. A significance of around 3 for $p_T > 4 \text{ GeV}/c$ is currently projected with $\mathcal{L}_{\text{int}} = 18 \text{ fb}^{-1}$ of pp collisions. The multi-charm baryon, such as Ξ_{cc}^{++} can be observed in Run 5, a novel detector with a high read-out rate, excellent pointing resolution, and outstanding tracking and particle identification over a large acceptance using advanced silicon detectors. The new area of charm hadronization will be accessed with the measurements of multi-charm baryons, such as Ξ_{cc}^{++} .

Monte Carlo simulation study for the medium response of bottomonia in p-A and A+A collisions is described in Chapter 6. To characterize the initial condition and evolution of the collision, the public MC framework, such as MG-Glauber and SONIC, are used, and the simulation framework has been developed based on the theoretically calculated thermal width of the $\Upsilon(nS)$ [219]. Only the dissociation effect is considered for calculating small collision systems, pp, p-Pb, and O-O colli-

sions at $\sqrt{s_{\text{NN}}} = 8$ TeV. The feed-down contributions from higher excited states are taken into account. The nuclear modification factor and elliptic flow of $\Upsilon(nS)$ are evaluated and compared to experimental measurements in Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The results from the SHINCHON show good agreement with the measurements. The sequential suppression for $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ is shown in the nuclear modification factor as a function of the number of participants in all small collision systems. In all small collision systems, the model predicts $v_2 < 0.01$ for $\Upsilon(nS)$, which is also consistent with the measurements in p-Pb collisions [43, 44]. The results from the SHINCHON can suggest the reference for the upcoming LHC run with the oxygen ion and can offer useful information on nuclear effects on the production of bottomonia in small systems.



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