

Study of K^{*0} and ϕ meson Production in proton+proton Collisions with ALICE at the LHC and Application of Non-extensive Statistics in High Energy Physics

Ph.D. Thesis

By

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**Study of K^{*0} and ϕ meson Production in
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Non-extensive Statistics in High Energy
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Arvind Khuntia



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CANDIDATE'S DECLARATION

I hereby certify that the work which is being presented in the thesis entitled **"Study of K^0 and ϕ meson Production in proton+proton Collisions with ALICE at the LHC and Application of Non-extensive Statistics in High Energy Physics"** in the partial fulfilment of the requirements for the award of the degree of **DOCTOR OF PHILOSOPHY** and submitted in the **DISCIPLINE OF PHYSICS, Indian Institute of Technology Indore**, is an authentic record of my own work carried out during the time period from July 2014 to December 2018 under the supervision of Prof. Raghunath Sahoo, Associate Professor, Discipline of Physics, Indian Institute of Technology Indore, Indore, M.P., India.

The matter presented in this thesis has not been submitted by me for the award of any other degree of this or any other institute.

Arvind Khuntia
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This is to certify that the above statement made by the candidate is correct to the best of my knowledge.

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Dedicated
to
My Family

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(Arvind Khuntia)

ABSTRACT

In today's universe, all ordinary matter is made up of atoms and these atoms are made up of protons and neutrons (except hydrogen) surrounded by cloud of electrons. The protons and neutrons are made up of quarks and bound together by gluons. The strong interactions between quarks are described by the theory of Quantum Chromo Dynamics (QCD). The QCD has two important properties because of the non-abelian nature of the gauge bosons. At high energy or momentum transfer (Q^2) the strong coupling constant becomes so small as if the interacting partons are free which is known as asymptotic freedom. On the other hand, at lower-momentum transfer, the coupling constant becomes high and hence it becomes impossible to separate the individual quarks to observe them free. This is because of the confining nature of strong interaction. These properties can lead to interesting phases of matter depending upon the temperature and/or density as the quarks and gluons become the relevant degrees of freedom instead of hadrons under such extreme circumstances. It has been speculated that a few micro-seconds after the Big-bang when the temperature was 100 MeV (10^{12} K) the universe was filled with quarks and gluons which are not free but deconfined. This phase is known as quark-gluon plasma (QGP). To create and study the QGP phase is the main motivation behind the ultra-relativistic heavy-ion collision experiments, where extreme temperatures/densities can be achieved. The experimental search for the de-confined state of quarks and gluons got a boost from the year 2000 at the Relativistic Heavy Ion Collider (RHIC) at BNL, where Au ions collide at $\sqrt{s_{NN}} = 20 - 200$ GeV and thereafter in the Large Hadron Collider (LHC) at CERN. All four RHIC experiments (BRAHMS, PHOBOS, PHENIX, and STAR) have discovered a new phase of nuclear matter that exhibits the properties of the strongly interacting, nearly perfect liquid.

At RHIC, various signatures such as jet quenching, azimuthal anisotropy, J/ψ suppression, and strangeness enhancement indicated a creation of a deconfined state of quarks and gluons known as QGP. A Large Ion Collider Experiment (ALICE) at the LHC is specifically designed for the creation of QGP and study its properties. The system formed in the heavy ion collisions undergoes a phase transition from a state of quark-gluon plasma to a phase of hadron gas. During this phase transition, the boundary at which the inelastic processes cease known as the chemical freeze-out and similarly at kinetic freeze-out boundary the elastic processes cease. The phase between chemical freeze-out and kinetic freeze-out boundaries is known as the hadronic phase. The lifetime of the fireball created in high energy collisions is comparable to the lifetime of resonance particles such as $K^*(892)^0$ and $\phi(1020)$, which have lifetimes $4.16 \pm 0.05 \text{ fm}/c$ and $46.3 \pm 0.4 \text{ fm}/c$. By studying the resonance particle we can probe the hadronic phase produced in high energy hadronic and nuclear collisions.

The $\phi(1020)$ meson is one of the lightest vector mesons consisting of $(s\bar{s})$, helps in the study of the strangeness production. Similarly, $K^*(892)^0$ is also a vector meson having a similar mass to $\phi(1020)$ but differs in strangeness content by one unit. A combined study of both the mesons helps to probe the strangeness production mechanism. Here, we study the minimum bias measurements of $K^*(892)^0$ and $\phi(1020)$. These results include: the transverse momentum spectra (p_T), mean transverse momentum, differential yields and also particle ratios with respect to stable particles. These results also provide an input to tune the perturbative-Quantum Chromodynamics (pQCD) inspired Monte Carlo (MC) event generators such as PYTHIA, PHOJET, and EPOS. Furthermore, the measurements in minimum bias pp collisions at $\sqrt{s} = 8 \text{ TeV}$ reported in this thesis serve as reference data to study nuclear effects in proton-lead (p-Pb) and Pb-Pb collisions.

High multiplicity pp-collisions at the LHC give us a possible opportunity to study matter under extreme conditions *i.e.* at high temperature and/or energy density and hence to explore the possibility of a formation of a QGP droplet in pp-collisions. The kinetic freeze-out surface of identified particles are determined by studying the transverse momentum spectra (p_T) of the produced par-

ticles in high energy collisions. Here, we analyse the transverse momentum (p_T)-spectra as a function of charged-particle multiplicity at midrapidity ($|y| < 0.5$) for various identified particles in proton-proton collisions at $\sqrt{s} = 7$ TeV using Boltzmann-Gibbs Blast Wave (BGBW) model and thermodynamically consistent Tsallis distribution function. We obtain the multiplicity dependent kinetic freeze-out temperature (T_{kin}) and radial flow (β) of various particles after fitting the p_T -distribution with BGBW model. Here, T_{kin} exhibits mild dependence on multiplicity class while β shows almost independent behaviour. The information regarding Tsallis temperature and the non-extensivity parameter (q) are drawn by fitting the p_T -spectra with Tsallis distribution function. The extracted parameters of these particles are studied as a function of charged particle multiplicity density ($dN_{ch}/d\eta$). In addition to this, we also study these parameters as a function of particle mass to observe any possible mass ordering. All the identified hadrons show a mass ordering in temperature, non-extensive parameter and also a strong dependence on multiplicity classes, except the lighter particles. It is observed that as the particle multiplicity increases, the q -parameter approaches to Boltzmann-Gibbs value ($q \sim 1$) indicating the tendency of the system to approach a thermal equilibrium. The observations are consistent with a differential freeze-out scenario of the produced particles. To explain the identified particle spectra in heavy-ion collisions, incorporation of radial flow in Tsallis statistics is necessary and thus assuming $(q - 1)$ close to zero, a Taylor series expansion of Tsallis distribution function in powers of $(q - 1)$ is provided. This helps in studying the degree of deviation of transverse momentum spectra and other thermodynamic quantities from a thermalized Boltzmann distribution. After checking thermodynamic consistency, we provide analytical results for the Tsallis distribution in the presence of collective flow up to first order in $(q - 1)$. The theoretical expectations are compared with the experimental data. Similarly, the speed of sound (c_s) is studied to understand the hydrodynamical evolution of the matter created in heavy-ion collisions. The quark-gluon plasma (QGP) formed in heavy-ion collisions evolves from an initial QGP to the hadronic phase via a possible mixed phase. Due to the system expansion in a first order phase transition scenario, the speed of sound reduces to zero as the specific heat diverges. We study the

speed of sound for systems, which deviate from a thermalized Boltzmann distribution using non-extensive Tsallis statistics. In the present work, we calculate the speed of sound as a function of temperature for different q -values for a hadron resonance gas. We observe a similar mass cut-off behaviour in the non-extensive case for c_s^2 by including heavier particles, as is observed in the case of a hadron resonance gas following equilibrium statistics. Also, we explicitly show that the temperature where the mass cut-off starts to vary with the q -parameter which hints at a relation between the degree of non-equilibrium and the limiting temperature of the system. It is shown that for values of q above approximately 1.13, all criticality disappears in the speed of sound, i.e. the decrease in the value of the speed of sound, observed at lower values of q , disappears completely.

Again, thermodynamic observables such as specific heat and isothermal compressibility are very useful in quantifying the nature of the phase transition and to obtain the equation of state (EOS) of the produced matter. Here, we have studied the isothermal compressibility (κ_T) as a function of temperature, baryon chemical potential and centre-of-mass energy ($\sqrt{s_{NN}}$) using hadron resonance gas (HRG) and excluded-volume hadron resonance gas (EV-HRG) models. A mass cut-off dependence of isothermal compressibility has been studied for a physical resonance gas. Further, we study the effect of heavier resonances (> 2 GeV) on the isothermal compressibility by considering the Hagedorn mass spectrum, $\rho(m) \sim \exp(bm)/(m^2 + m_0^2)^{5/4}$. Here, the parameters, b and m_0 are extracted after comparing the results of recent lattice QCD simulations at finite baryonic chemical potential. We find a significant difference between the results obtained in EV-HRG and HRG models at higher temperatures. The inclusion of the Hagedorn mass spectrum in the partition function for hadron gas has a large effect at a higher temperature. A higher mass cut-off in the Hagedorn mass spectrum takes the isothermal compressibility to a minimum value, which occurs near the Hagedorn temperature (T_H). Furthermore, it is always exciting to explore possible applications of non-extensivity in condensed matter physics- a cross-area application. Here, we use Tsallis thermodynamically consistent distribution in studying the thermal effective mass and specific heat of various metals. We have found an excellent agreement with experimental findings. These works form the

basis of the present thesis.

The thesis is organised as follows: after a motivation to study pp collisions at the LHC energies in contrast to the QGP like behaviour in heavy-ion collisions, a brief discussion on the QCD phase diagram and possible phase transitions, and location of CP is discussed in Chapter 1. In Chapter 2, we discuss about the ALICE detectors used for our data analysis with more focus to ITS, TPC, TOF and V0 detectors. In Chapter 3, we present the study of K^{*0} and ϕ mesons production in pp-collisions with the ALICE detector at the LHC. The application of Tsallis non-extensive statistics and BGBW model in high-energy physics is provided in Chapter 4; along with the important findings of potential application in condensed matter physics as well. At the end, in Chapter 5 the results are summarised with important findings.

PUBLICATIONS

List of Publications:

(Results from papers marked with [*] are included in this thesis.)

1. [*] “Multiplicity Dependence of Non-extensive Parameters for Strange and Multi-Strange Particles in Proton-Proton Collisions at $\sqrt{s} = 7$ TeV at the LHC”
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2. [*] “Speed of sound in hadronic matter using non-extensive Tsallis statistics”
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3. [*] “Radial Flow in Non-Extensive Thermodynamics and Study of Particle Spectra at LHC in the Limit of Small $(q - 1)$ ”
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5. [*] “Effect of Hagedorn States on Isothermal Compressibility of Hadronic Matter formed in Heavy-Ion Collisions: From NICA to

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6. [*] “ **$K^*(892)^0$ and $\phi(1020)$ meson production at high transverse momentum in pp and Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV**”

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9. [*] “**Measurement of ϕ Mesons at Mid-Rapidity in Minimum-Bias pp Collisions at $\sqrt{s} = 8.0$ TeV**”

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10. [*] “ **$K^*(892)^0$ and $\phi(1020)$ production in minimum bias pp collisions at $\sqrt{s} = 8$ TeV**”

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12. **“Transverse Momentum Spectra and Nuclear Modification Factor using Boltzmann Transport Equation with Flow in Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV”**

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13. **“Inelastic Normalization Factors for pp collisions at $\sqrt{s} = 5.02, 8,$ and 13 TeV”**

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14. **Additional 85 publications in ALICE experiment**

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2. [*] “The q-Statistics and QCD Thermodynamics at LHC”
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3. [*] “Speed of Sound in a System Approaching Thermodynamic Equilibrium”
A. Khuntia, P. Sahoo, P. Garg, R. Sahoo and J. Cleymans.
DAE Symp. Nucl. Phys. **61**, 842 (2016).
4. [*] “Radial Flow in Non-Extensive Thermodynamics and Study of Particle Spectra at LHC in the Limit of Small (q-1)”
T. Bhattacharyya, J. Cleymans, A. Khuntia, P. Pareek and R. Sahoo.
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5. “Systematic Study of Identified Particle Spectra with Radial Flow at RHIC and LHC in Tsallis Non-extensive Statistics”
D. Thakur, P. Garg, S. Tripathy, A. Khuntia, A. N. Mishra and R. Sahoo.
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Chapter 1

Introduction

1.1 The Standard Model

In today's Universe, the constituents of the ordinary matter are atoms and these are made up of nucleus (consists of protons and neutrons (except hydrogen atom)) surrounded by cloud of electrons. For a long time, it was believed that 'Atomic Nucleus' supposed to be composed of protons and neutrons and they strongly interact with each other. However, later Murray Gell-Mann and George Zweig (1964) found that all these particles, including protons and neutrons, consist of only a few building blocks which he termed as quarks. The quarks and leptons don't have substructures and called as elementary/fundamental particles. There are four fundamental forces, (i) strong (ii) electromagnetic (iii) weak and (iv) gravitational in nature. Understanding of how these particles and first three of the forces are related to each other is encapsulated in the Standard Model of particle physics by Glashow, Salam, and Weinberg in 1970s. According to the Standard model, quarks and leptons are classified into three generations as shown in Fig. 1.1. The up (u) and down (d) quarks, the charm (c) and strange (s) quarks, and, top (t) and bottom (b) quarks belong to first, second and third generations in the standard model, respectively. Similarly, the leptons are categorized into three generations; electron, muon, tau and their corresponding neutrinos. The quarks and leptons are fermions and have a half integral spin quantum number.

First three of the fundamental forces result from the exchange of force carrier particles known as mediators (bosons) and these are experimentally discovered.

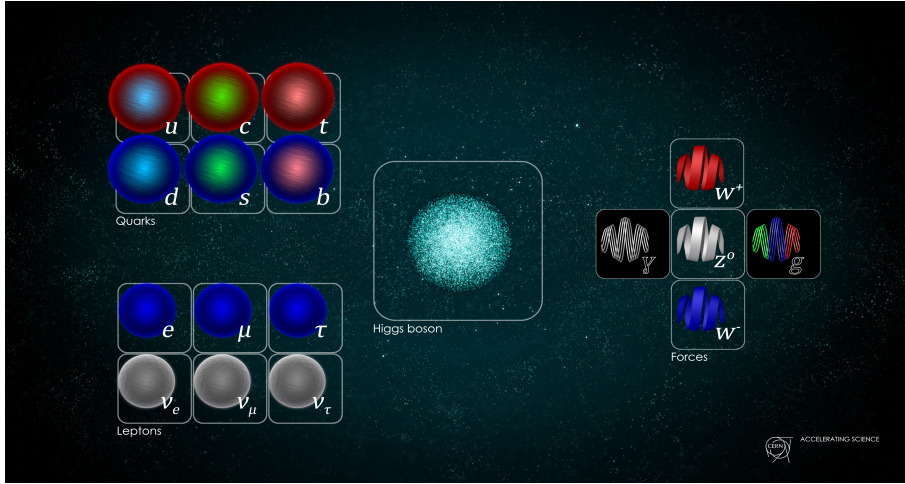


Figure 1.1: (color online) Three generation of matter particles in the standard model. The mediator particles (Gauge bosons) are also shown [1].

A discrete amount of energy is transferred among these particles by exchanging bosons with each other. Each fundamental force has its own force carrier; the strong force is carried by the “gluons”, the electromagnetic force is carried by the “photons”, and the “W and Z bosons” are responsible for the weak force. The fourth force carrier particle corresponding to gravity is the graviton, which is not yet discovered. The Standard Model includes the electromagnetic, strong and weak forces and their carrier particles. However the fourth force, which we experience in our daily lives, i.e, gravity is not included in the standard model as fitting gravity comfortably into this framework has proved to be a difficult challenge and still it is an open problem. The theory of quantum mechanics used to describe the micro-world, and the general theory of relativity used to describe the macro-world, are difficult to fit into a single framework. No one has managed to make the two mathematically compatible in the context of the Standard Model. For particle physics, when it comes to the minuscule scale of particles, gravity is so weak that we can neglect its effect. Gravity has a significant effect when the matter is in bulk, at the scale of the human body or of the planets. That’s why the Standard Model still works well despite its reluctant exclusion of one of the fundamental forces. Even though the Standard Model is

the best theory for the description of the micro world, it explains only three out of four fundamental forces. Also, it doesn't explain the existence of dark matter, the matter-antimatter asymmetry in our Universe and also the finite mass of the neutrinos, which are the main evidences of the existence of physics beyond the Standard Model.

1.2 Quantum Chromodynamics

The strong interactions between quarks are described by the theory of Quantum Chromo Dynamics (QCD). The quarks and gluons have colour charge but they constitute to form colourless hadrons. The static quark-antiquark potential which takes into account both the long-range and short-range behaviour of QCD interaction is given as

$$V_{QCD}(r) = -\frac{4}{3} \frac{\alpha_S}{r} + kr, \quad (1.1)$$

where α_S is the strong coupling constant or known as running coupling constant, k is colour string tension constant and r is the distance between interacting partons. The QCD running coupling constant α_S is given as,

$$\alpha_S(Q^2) = \frac{12\pi}{(33 - 2n_f) \ln(Q^2/\Lambda_{QCD}^2)}, \quad (1.2)$$

where Q^2 is momentum transfer between partons, n_f is the number of quark flavours and $\Lambda_{QCD} \sim 200$ MeV is the QCD scale parameter. Looking at the strong potential at lower-momentum transfer, the coupling constant becomes high and hence it becomes impossible to separate the individual quarks to observe them as free. This is because of the confining nature of strong interaction. On other hands, at high energy or momentum transfer (Q) the strong coupling constant becomes so small as if the interacting partons are free which is known as asymptotic freedom. The behaviour of strong running coupling constant (α_S) with respect to momentum transfer (Q) and the values estimated by different experiments are shown in Figure 1.2.

These properties of QCD can lead to interesting phases of matter depending upon the temperature and/or density as the quarks and gluons become the

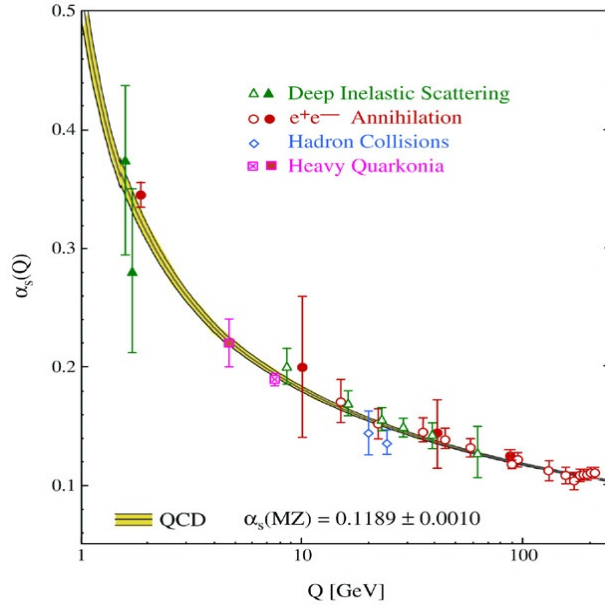


Figure 1.2: QCD running coupling constant (α_s) as a function of momentum transfer (Q) [2, 3].

relevant degrees of freedom instead of hadrons under such extreme circumstances. It has been speculated that a few microseconds after the Big-bang when the temperature was 100 MeV (10^{12} K) the universe was filled with quarks and gluons which are not free but deconfined. This phase is known as quark-gluon plasma (QGP).

1.3 Quark Gluon Plasma

To recreate conditions similar to those of the very early Universe, powerful accelerators make head-on collisions of massive ions and hundreds of protons and neutrons are smashed into one another at ultra-relativistic energies of a few trillion electron-volts each.

Lattice-QCD (lQCD) predicts the QCD phase transition from a confined phase of hadron to deconfined partonic phase at lower baryon chemical potential (μ_B) [4]. The lQCD calculation considering massless quarks at zero chemical potential ($\mu_B = 0$) and at high-temperature limit, the temperature (T) dependence of the energy density (ε) scaled by T^4 , which also tells about the degrees of freedom is shown in Fig 1.3.

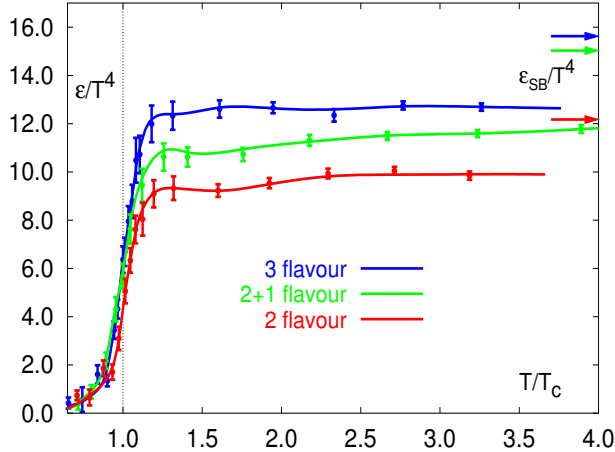


Figure 1.3: lQCD predictions of ϵ/T^4 as a function of temperature normalized by the critical temperature (T_C). The horizontal arrows show the Stefan-Boltzmann limit [5].

The scaled energy density (ϵ/T^4) shows a sharp rise near the T_c ($T_C \sim 154$ TeV (for 3 – flavour QCD)) and it saturates at higher temperatures below the Stefan-Boltzmann limit [5]. This behaviour shows an abrupt rise in degrees of freedom of the system and indicates a change of phase from hadron gas to a deconfined state of quarks-gluon plasma. This hints for a critical energy density and a critical temperature for the QCD phase transition.

1.3.1 The QCD phase diagram

The conjectured QCD phase diagram of strongly interacting matter for different phases of QCD matter and possible phase transition is shown in Fig. 1.4.

Here, the state of the strongly interacting matter is shown with the variation of temperature (T) and the baryon chemical potential μ_B . While studying the particle yields at the chemical freeze-out in the heavy-ion collisions [6] and their comparison to statistical models suggests that with center-of-mass energy ($\sqrt{s_{NN}}$), T increases and μ_B decreases. Hence by varying the center-of-mass energy ($\sqrt{s_{NN}}$), we can get access to a large part of the phase space. The phase diagram shows a rich phase structure, however, experimentally only some part is accessible which corresponds to some of the following distinct structures: de-

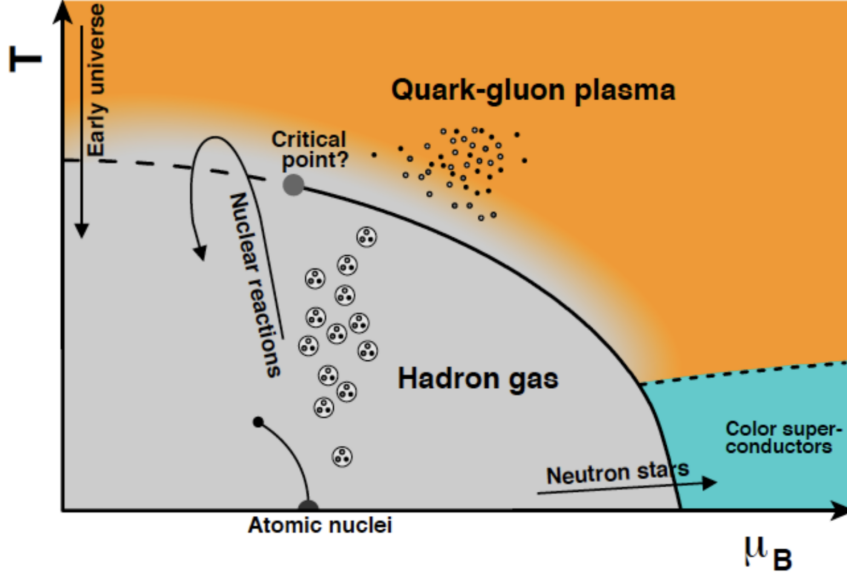


Figure 1.4: (Color online) Conjectured phase diagram of strongly interacting matter.

confined phase of quarks and gluons, hadronic phase, critical point and a crossover at low μ_B and very high temperature.

1.3.2 Ultra-relativistic collisions

The experimental search for the de-confined state of quarks and gluons got a boost from the year 2000 at the Relativistic Heavy Ion Collider (RHIC) at BNL, where Au ions collide at $\sqrt{s_{NN}} = 20 - 200$ GeV and thereafter in the Large Hadron Collider (LHC) at CERN. All four RHIC experiments (BRAHMS, PHOBOS, PHENIX, and STAR) have discovered a new phase of nuclear matter that exhibits the properties of the strongly interacting, nearly perfect liquid. At RHIC, various signatures such as jet quenching, azimuthal anisotropy, J/ψ suppression, and strangeness enhancement indicated a creation of a deconfined state of quarks and gluons known as QGP. A Large Ion Collider Experiment (ALICE) at the LHC is specifically designed for the creation of QGP and study its properties. In the high-energy experiments, two nuclei are accelerated to ultra-relativistic velocities resulting in the collision of two Lorentz contracted nuclei along the direction of motion. The overlapping area at the interaction point depends on the impact

parameter of the colliding nuclei.

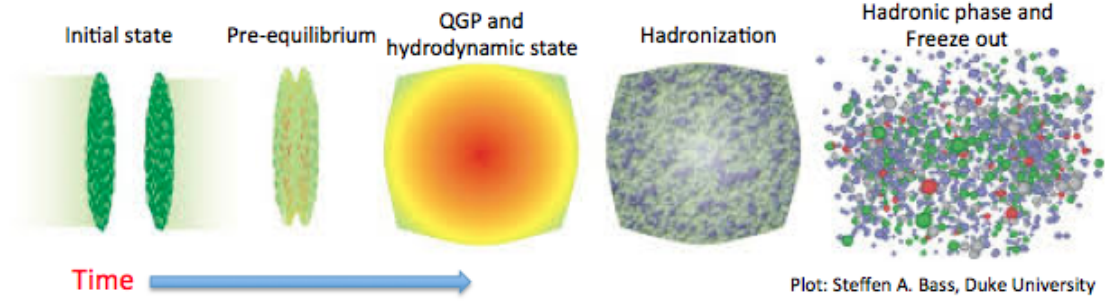


Figure 1.5: (Color online) A schematic picture of the time evolution of heavy-ion collision. (Image credit: Prof. Steffen A. Bass).

The system formed in the heavy ion collisions undergoes a phase transition from a state of quark-gluon plasma to a phase of hadron gas through different stages as shown in Fig.1.5 and 1.6 .

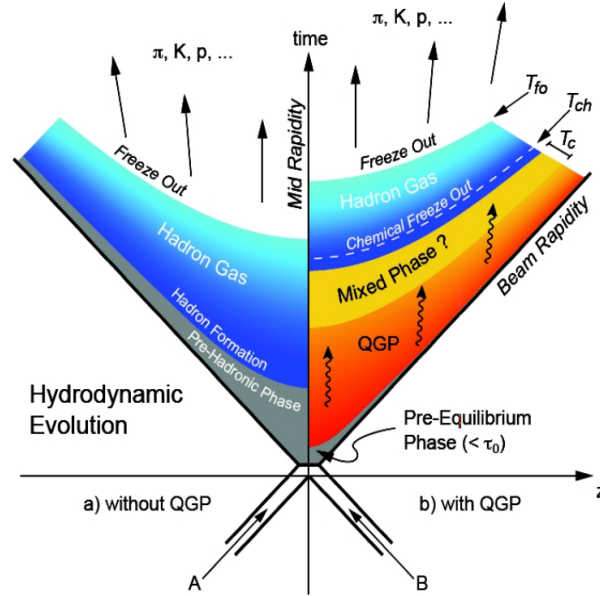


Figure 1.6: (Color online) A schematic diagram of the longitudinal space-time evolution of relativistic nucleus-nucleus collision [7].

During this phase transition, the boundary at which the inelastic processes cease known as the chemical freeze-out and similarly at kinetic freeze-out boundary the elastic processes cease. The phase between chemical freeze-out and kinetic freeze-out is known as the hadronic phase. The produced deconfined QGP state

only exists for a few microseconds and it is impossible to observe this state within such a small lifetime. However, there are various indirect observables which can be used as diagnostic tools for observation of QGP. There is no unique signal which alone gives confirmation for the production of a QGP phase. Therefore, a comprehensive understanding of the signatures of the QGP is necessary. Various QGP signatures like azimuthal anisotropy, strangeness enhancement, and J/ψ suppression are discussed below in details.

- **Elliptic flow and QGP:**

Elliptic flow is one of the important observables to understand the collective behaviour of the system produced in heavy-ion collisions and is also sensitive to the initial geometry of the colliding nuclei. Due to the initial spatial anisotropy of the produced system, the pressure gradient is different in different directions, which results in the final state momentum anisotropy. The azimuthal momentum anisotropy is defined as the Fourier expansion in the azimuthal angle ϕ as [8, 9],

$$\frac{dN}{p_T dp_T d\phi dy} = \frac{dN}{2\pi p_T dp_T d\phi dy} \left[1 + 2 \sum_1^{\infty} v_n \cos[n(\phi - \psi)] \right]. \quad (1.3)$$

Where v_n 's are various order of flow harmonics, v_1 is the directed flow, v_2 is the elliptic flow, v_3 is triangular flow and so on. ψ is reaction plane angle. Reaction plane is the plane formed by the beam direction (usually the longitudinal direction i.e. along z-axis) and impact parameter vector (transverse distance between the centre of two nuclei).

Fig. 1.7 shows that the elliptic flow of the dense media produced at RHIC and LHC are of similar nature. The 30% higher value of v_2 at LHC is due to the availability of larger phase space which increases the yield of high- p_T particles at LHC. The large elliptic flow suggests larger rescattering among partons. This also indicates an early thermalization of high- p_T partons. All these above evidences signify a strong signature for a formation of QGP. The v_2 scaled with the number of constituent quarks (NCQ) as a function of $(m_T - m_0)/n_q$, where $m_T = \sqrt{p_T^2 + m_0^2}$ and m_0 is mass of identified particles, shown in right panel of Fig. 1.7. The observation of number of

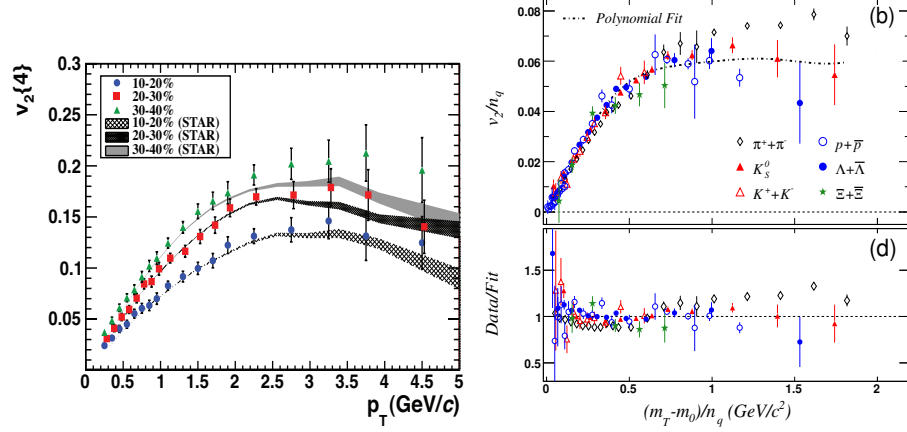


Figure 1.7: Left: The p_T differential elliptic flow of charged particles. The Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV shown in coloured symbols are compared to Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV, which are in shaded regions [10]. Right: v_2/n_q versus $(m_T - m_0)/n_q$ of identified hadrons in Au–Au collisions at $\sqrt{s_{NN}} = 62.4$ GeV [11].

constituent quark scaling suggests hadron formation via quark coalescence. The NCQ scaling of v_2 also indicates that the relevant degrees of freedom in the early system may be sub-hadronic (e.g. constituent quarks).

- **J/ ψ Suppression:**

J/ ψ is a bound state of charm quark and its anti-quark, $c\bar{c}$, usually formed at the early stages of the hadronic and nuclear collisions due to its heavier mass. That's why the J/ ψ becomes an excellent probe to study the early stage dynamics of the medium formed in hadronic/heavy-ion collisions. The production of J/ ψ in heavy-ion collisions, where a QGP state is expected will be suppressed as compared to pp collisions as suggested by theoretical calculations [12]. Usually no medium formation is expected in pp-collisions [13].

This J/ ψ suppression was observed up to top RHIC energies. At LHC energies, the dynamics of the particle production changes bringing the interplay between the regeneration process and suppression due to Debye-like screening. The QGP medium screen the colour charge of the $c\bar{c}$, which reduces the production of J/ ψ .

The J/ψ suppression is clearly observed in heavy-ion collisions and the disassociated charm quarks contribute towards enhancement of open charm production ($D^0, D^\pm$). The J/ψ suppression is first observed in SPS [14] and then in RHIC [15, 16] energies confirming to formation of a QGP medium. The nuclear modification factor (R_{AA}) of J/ψ meson in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV and 5.02 TeV as a function of centrality measured in ALICE at LHC is shown in Fig. 1.8. The results of Au+Au collisions at $\sqrt{s_{NN}} = 0.2$ TeV from the PHENIX collaboration at RHIC is also shown. The suppression at LHC energies is found to be less than that of RHIC and SPS energies [15]. This enhancement in the J/ψ yield at LHC is explained by the recombination of charm quarks during the hadronization process. So the suppression in R_{AA} at $\sqrt{s_{NN}} = 5.02$ TeV is less compared to $\sqrt{s_{NN}} = 2.76$ and 0.2 TeV.

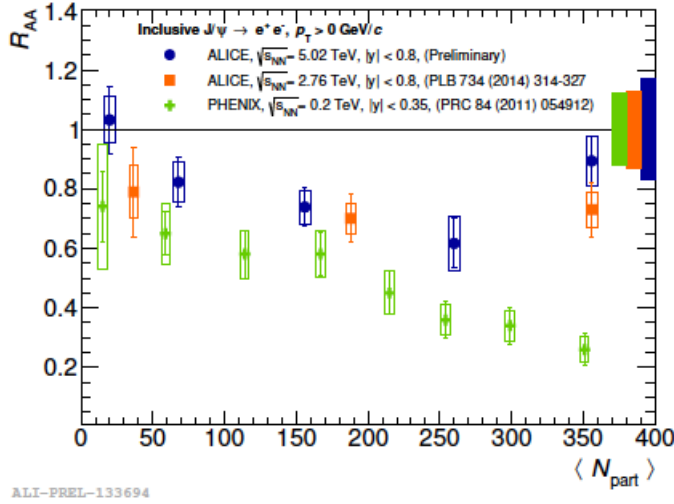


Figure 1.8: (Color online) The R_{AA} of J/ψ as function of centrality in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV and 5.02 TeV at ALICE and the measurement is compared with the PHENIX result of Au+Au collisions at $\sqrt{s_{NN}} = 0.2$ TeV.

- **Strangness Enhancement:**

The production of strange hadrons in high-energy collisions provides a unique tool to study the strongly interacting matter, which is governed by partonic degrees of freedom and is described by Quantum Chromo Dynamics (QCD). Considering the constituent quarks, in hadronic and nuclear

collisions, only the light quarks, namely u and d , participate as the colliding quanta of the system. The strange quarks are produced via hard processes, by annihilation of light quark-antiquark pairs ($q\bar{q} \rightarrow s\bar{s}$) and gluon fusion ($gg \rightarrow s\bar{s}$) [17] during the early stages of the collisions. The enhancement of the strange hadron production thus serves as one of the signatures of Quark-Gluon Plasma (QGP), where quarks and gluons are no longer confined within the hadronic domain, rather the partonic degrees of freedom become important in nuclear dimensions. The strange particle production mechanism and production rates are different both in hadron gas and QGP media. Strangeness enhancement can be quantified by studying the abundances of the strangeness between plasma and hadronic phase.

In a high gluon density environment in QGP medium, the formation of $s\bar{s}$ pair from the channel, $gg \rightarrow s\bar{s}$ dominates over the annihilation of light quarks. In other hand, the annihilation of light quarks to strange quark is the main channel for the production of strangeness in pp collisions. This distinguishes a QGP phase from a hadronic phase.

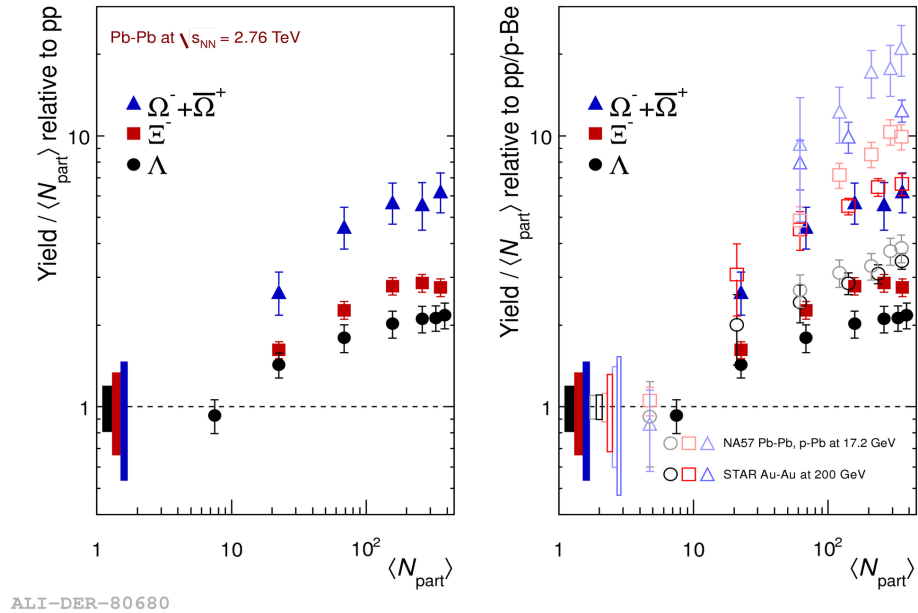


Figure 1.9: The yield of multi-strange hadrons in Pb-Pb relative to pp collisions measured in ALICE (left panel) and NA57, STAR (right panel) as function of centrality, $\langle N_{\text{part}} \rangle$) [19–21].

The enhancement factor E is defined as,

$$E = \frac{2}{\langle N_{part} \rangle} \left[\frac{\frac{dN^{AA}}{dy}|_{y=0}}{\frac{dN^{pp}}{dy}|_{y=0}} \right]. \quad (1.4)$$

Here $\langle N_{part} \rangle$ is the number of participant nucleons, which is a measure of collision centrality. $\frac{dN^{AA}}{dy}|_{y=0}$ and $\frac{dN^{pp}}{dy}|_{y=0}$ are the yields in heavy-ion and pp collisions, respectively.

The yield of hyperons enhancement factor (E) as a function of centrality N_{part} is shown In Fig. 1.9 for ALICE [19] and is compared with measurements from STAR and NA57 collaborations [20, 21]. The observed abundant production of strange particle in Pb–Pb with respect to pp and p–Be collisions indicates the formation of deconfined state of matter at LHC energies. It is also found that the enhancement factor is higher for particles containing more strange quark i.e. $E(\Lambda) < E(\Xi) < E(\Omega)$ [19–21].

1.4 The Role of pp-collisions in Understanding a QGP Phase and the Basis of this Thesis

Proton-proton collisions at a given center-of-mass energies serves as a baseline to study a possible QGP phase and nuclear effects in heavy-ion collisions. In this thesis work, we have attempted to study pp-collisions at LHC energies in focus with resonances. The dynamics of resonance production can be used as a powerful and important tool to understand the QGP medium formed at the LHC, which is one of the main motivations.

The lifetime of the fireball created in high energy collisions is comparable to the lifetime of resonance particles such as $K^{*0}(892)$ and $\phi(1020)$, which have lifetimes $4.16 \pm 0.05 \text{ fm}/c$ and $46.3 \pm 0.4 \text{ fm}/c$. By studying the resonance particles, we can probe the hadronic phase produced in high energy hadronic and nuclear collisions. The $\phi(1020)$ is one of the lightest vector mesons consisting of $(s\bar{s})$, which helps in the study of the strangeness production. Similarly, $K^{*0}(892)$ is also a vector meson having a similar mass to $\phi(1020)$ but differs in strangeness content by one unit which also helps to probe the strangeness production. The measurement of particle production in minimum bias pp collisions

provides an input to tune the perturbative-Quantum Chromodynamics (pQCD) inspired Monte Carlo (MC) event generators such as PYTHIA, PHOJET, and EPOS. Furthermore, the measurements in minimum bias pp collisions at $\sqrt{s} = 8$ TeV reported in this thesis serve as reference data to study nuclear effects in proton-lead (p-Pb) and Pb-Pb collisions.

High multiplicity pp-collisions at LHC give us the opportunity to study matter under extreme conditions *i.e.* at high temperature and/or energy density. The kinetic freeze-out surface of identified particles can be determined by studying the transverse momentum spectra (p_T) of the produced particles in high energy collisions. Higher probability of multi-partonic interactions at higher collision energies leads to multiple interactions in the produced system, which results in high-multiplicity pp-collisions. These multi-partonic interactions might lead to thermalization in high multiplicity pp-collisions, which can be described by statistical models. The non-exponential behaviour of the transverse momentum spectra at high- p_T is not explained by the Boltzmann-Gibbs distribution function. This can be achieved by considering the thermodynamically consistent Tsallis non-extensive statistics, which behaves like an exponential function at low- p_T and a power law function at high- p_T . The parameter q in the non-extensive statistics shows the deviation of the produced system from an equilibrium distribution function. Similarly, we can investigate any possible hints of collectivity by looking into the low- p_T part of the transverse spectra using the Boltzmann-Gibbs Blast Wave (BGBW) model. To explain the identified particle spectra in heavy-ion collisions, incorporation of radial flow in Tsallis statistics is necessary and thus assuming $(q - 1)$ close to zero, a Taylor series expansion of Tsallis distribution function in powers of $(q - 1)$ is provided. Similarly, the equation of state (EoS) of the system produced in high energy collisions can be understood by looking into the speed of sound in that medium. This also helps in looking for any critical behaviour the system could undergo during its space-time evolution. The study of the speed of sound in hadronic matter hence becomes a subject of paramount importance in heavy-ion collisions.

Again, Thermodynamic observables such as specific heat and isothermal compressibility are very useful in quantifying the nature of the phase transition

and to obtain the equation of state (EOS) of the matter. Specific heat measures the change in energy density with respect to the change in temperature and is also related to the temperature fluctuation in the system [22–24]. The isothermal compressibility (κ_T) defines the change in volume with the change in pressure at constant temperature. It can be expressed as a second order derivative of Gibbs free energy with respect to pressure and is expected to diverge at CP. Thus, κ_T and C_V are very helpful in unveiling the nature of the phase transition and helpful in the search for CP. Furthermore, it would be exciting to explore possible applications of non-extensivity in condensed matter physics. Here, we use Tsallis thermodynamically consistent distribution in studying the thermal effective mass and specific heat of various metals. We have found an excellent agreement with experimental findings. These works form the basis of the present thesis.

The thesis is organised as follows: after a motivation to study pp collisions at the LHC energies in contrast to the QGP like behaviour in heavy-ion collisions, a brief discussion on the QCD phase diagram and possible phase transitions, and location of CP is discussed. In Chapter 2, we discuss about the ALICE detectors used for our data analysis with more focus to ITS, TPC, TOF and V0 detectors. In Chapter 3, we present the study of K^{*0} and ϕ mesons production in pp-collisions with the ALICE detector at the LHC. The application of Tsallis non-extensive statistics and BGBW model in high-energy physics is provided in Chapter 4; along with the important findings of potential application in condensed matter physics as well. At the end, in Chapter 5 the results are summarised with important findings.

Chapter 2

A Large Ion Collider Experiment at the LHC

ALICE (A Large Ion Collider Experiment) [25], is one of the major detector located at the interaction point 2 along the Large Hadron Collider (LHC) ring, designed to study the physics of strongly interacting matter at high temperature or extreme energy density, where a deconfined state of quarks and gluons believe to form. The main focus of this chapter will be the description of the ALICE experiment. In Section 2.1, the LHC main features will be discussed along with the some parameters that is relevant for the proton beams which has been used in this analysis. Section 2.2 presents the overview of ALICE experimental setup. Finally, Section 2.3 mostly focuses on the detectors used for tracking and particle identification for analysis of K^{*0} and ϕ meson production in pp collisions at $\sqrt{s}=8$ TeV.

2.1 CERN Accelerator Complex

Figure 2.1 shows various experiments at LHC, at the European Organisation for Nuclear Research (CERN) along with the accelerator complex. The LHC is a two counter rotating ring of superconducting accelerator and collider with a circumference of 26.7 Km. There are four major experiments located in four different interaction points along the LHC ring. It was built to achieve the highest energy pp collisions in laboratory, with aim for the discovery of Higgs boson and hints

for physics beyond standard model. Indeed, the discovery of Higgs boson was announced on the 4th of July 2012 by the two LHC experiments, ATLAS and CMS. Similarly the third experiment, LHCb, is dedicated for precision measurements of heavy flavour decay. Finally the fourth experiment, ALICE, is dedicated for the physics of strongly interacting QCD matter through the heavy-ion physics program. The LHC is designed to achieve $\sqrt{s} = 14$ TeV in case of pp and $\sqrt{s_{NN}} = 5.5$ TeV for Pb-Pb collisions [26].

The accelerator complex at CERN is a series of machines sequentially arranged to accelerate particles to higher energies. The energy of beam of particles is boosted before injecting the beams into the next machine in the sequence. The proton beams are obtained by passing the hydrogen gas through an electric field, which strips off the electrons leaving with the protons to enter the accelerator. Each machine boosts the energy of beam of particles, before injecting the beam into the next machine. Protons are accelerated to the energy of 50 MeV by the Linear Accelerator 2, which is the first accelerator in the accelerator chain.

The beam from the LINAC 2 is then injected into the Proton Synchrotron Booster (PSB), which further boosts the energy of protons to 1.4 GeV and then Proton Synchrotron (PS) pushes the beam to 25 GeV. Finally before injecting the beams to LHC ring, the Protons are accelerated to 450 GeV by the Super Proton Synchrotron (SPS). The LHC ring consists of two beam pipes and the injected beam in one pipe circulates clockwise while the beam in the other pipe circulates anti-clockwise. To fill each LHC ring, it takes 4 minutes and 20 seconds and 20 minutes for the protons to reach their maximum energy of 6.5 TeV. The two beams are brought into collision inside four experiments, ALICE at point 2, ATLAS at point 1, CMS at point 5 and LHCb at point 8, where the total energy at the collision point is equal to 13 TeV.

Like energy, the second important parameter is the instantaneous luminosity, which is necessary for the exploration of rare events is the proportionality factor between the event rate (R) and the cross section of the process under study (σ_{int}):

$$R = L \sigma_{int}. \quad (2.1)$$

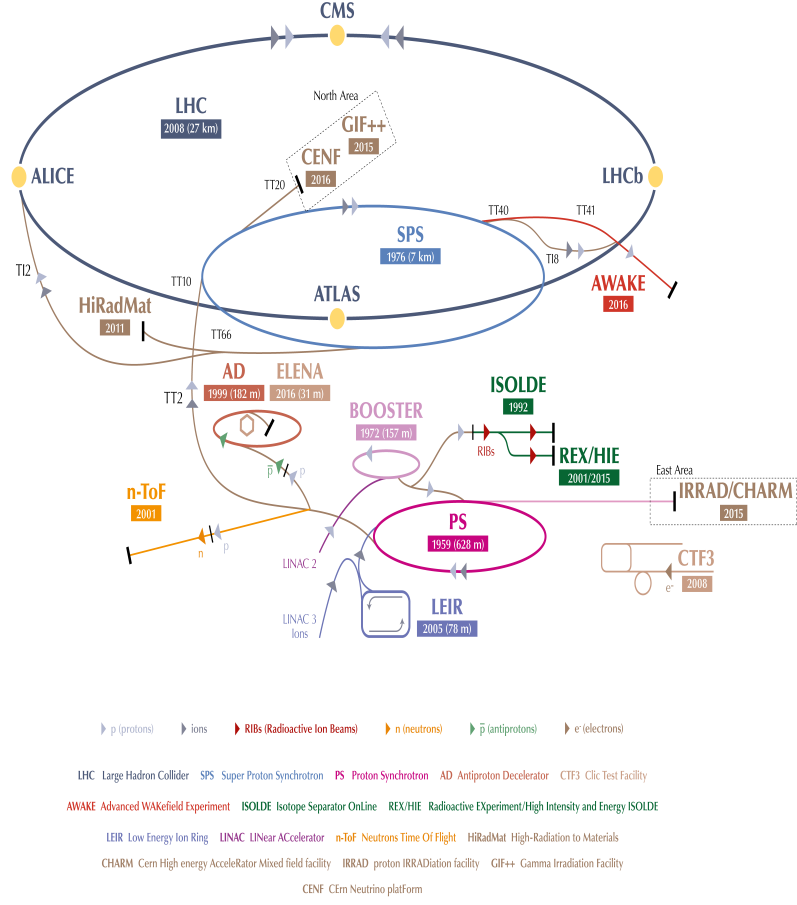


Figure 2.1: (Color online) The CERN accelerator complex [27].

The machine luminosity depends only on the colliding beam parameters at the interaction point and can be written for a gaussian beam distribution as given by,

$$L = \frac{N_b^2 n_b f_{rev} \gamma_r}{4 \pi \epsilon_n \beta^*} F \quad (2.2)$$

Where

N_b is the number of particles per bunch,

n_b the number of bunches per beam,

f_{rev} the revolution frequency,

γ_r the Lorentz gamma factor,

ϵ_n the normalised transverse emittance,

β^* the beta function at the collision point,

and F the geometrical reduction factor that accounts for the crossing angle at IP.

The transverse emittance and the beta function are related to the bunch cross-sectional size by,

$$\sigma_L^2 = \frac{1}{\pi} \beta \epsilon. \quad (2.3)$$

Lower transverse emittance of particle beam ensures that the particles are confined to a small distance and having nearly same momentum. Similarly the β -function is determined by the quadrupole focusing strength. Higher luminosity is achieved by keeping the transverse emittance low and β^* as small as possible. The transverse emittance changes with the beam momentum during acceleration hence the normalised emittance, $\epsilon_n = \gamma\beta\epsilon$ is often more useful to consider as it doesn't change during the acceleration.

The F factor can be obtained from the total crossing angle at the interaction point, θ_c , by assuming spatial distributions of particles in each bunch as Gaussian along the xy -plane as well as z direction and given by

$$F = \left[1 + \left(\frac{\theta_c \sigma_z}{2 \sigma_{xy}^*} \right)^2 \right]^{-1/2} \quad (2.4)$$

here σ_{xy} is the RMS beam size in the transverse plane and σ_z is the RMS bunch length. These bunch transverse and longitudinal sizes are same for both the bunches and the number of particles per bunch is $\simeq 10^{11}$. ATLAS and CMS are high luminosity pp experiments and run at a peak luminosity of $10^{34} \text{ cm}^{-2}\text{s}^{-1}$ while LHCb runs at the designed luminosity value of $10^{32} \text{ cm}^{-2}\text{s}^{-1}$. On the other hand the luminosity in ALICE experiment is limited to $10^{30} \text{ cm}^{-2}\text{s}^{-1}$ in order to avoid pile up in the main tracker detector.

During the first phase of LHC operations from the end of 2009 to the beginning of the technical shutdown in 2013 (Run 1) ALICE has accumulated huge amount of data at 0.9, 2.76, 7 and 8 TeV in pp collisions, p-Pb collisions at 5.02 TeV and Pb-Pb collisions at 2.76 TeV. In Run 2 ALICE carried out taking data with twice collision energy from 2015.

2.2 The ALICE Detector

ALICE is mainly aimed at studying the physics of strongly interacting QCD-matter produced in heavy-ion collisions. Thanks to the unique particle identification performance that allows a comprehensive study of the hadrons, electrons, muons and photons produced in the collision down to very low momentum (0.1 GeV) as well a good particle identification capabilities up to 20 GeV/ c . ALICE has been taking data for pp at both low and high energies, which will provide reference data for Pb–Pb. A full schematic view of Run 2 ALICE detector is shown in Fig. 2.2, which weights about 10,000-tonn with a dimension of 26 m long, 16 m high, and 16 m wide. A transverse schematic view of ALICE central barrel detectors are shown in Fig. 2.3.

ALICE uses a right-handed orthogonal Cartesian coordinate system. The origin of this coordinate system is at LHC interaction point 2 with z-axis parallel to the beam direction and points along the LHC Beam 2. The x-axis is horizontal and points towards the center of the LHC, similarly y-axis is approximately vertical and points upwards. ALICE is composed of seventeen detector system and is categorised into three groups

- Central-barrel detectors
- Forward detectors
- MUON spectrometer.

2.2.1 Central-barrel detectors

The central barrel detectors are group of detectors consisting of Inner Tracking System (ITS), Time Projection Chamber (TPC), Transition Radiation Detector (TRD), Time Of Flight (TOF), Photon Spectrometer (PHOS), Electromagnetic Calorimeter (EMCal), and High Momentum Particle Identification Detector (HMPID), which are placed in the L3 solenoid magnet with $B=0.5$ T. The first four detectors have full azimuthal acceptance at midrapidity ($|\eta| \leq 0.9$). The ITS and TPC are primary detectors for tracking of charged particles. The ITS is composed six layers of silicon detector, two Silicon Pixel Detectors (SPD), two

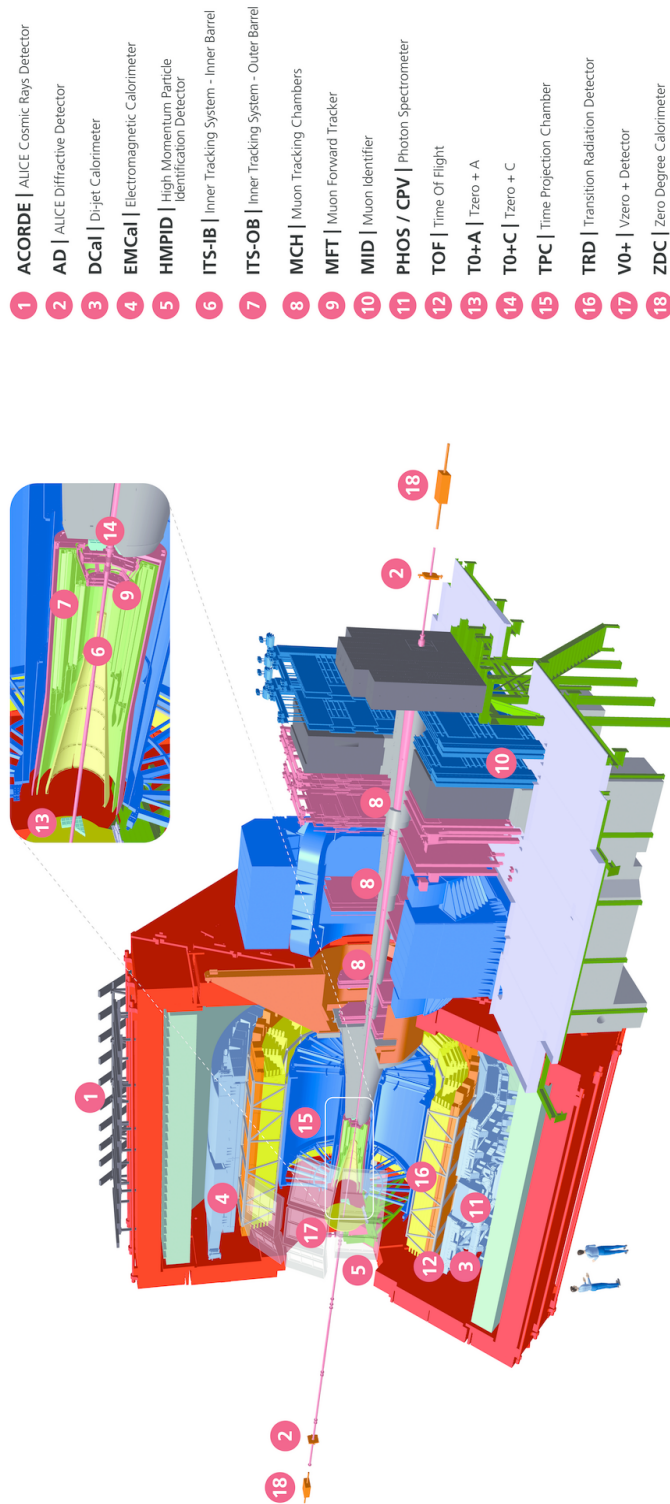


Figure 2.2: (Color online) The ALICE experiment at the CERN LHC [28].

Silicon Drift Detectors (SDD), and two Silicon Strip Detectors (SSD). TPC is a gaseous detector having 90 m^3 drift volume filled with Ne–CO₂ and is divided into two parts by the central cathode, which is kept at -100 kV. The end plates of the TPC are equipped with multiwire proportional chambers (MWPC).

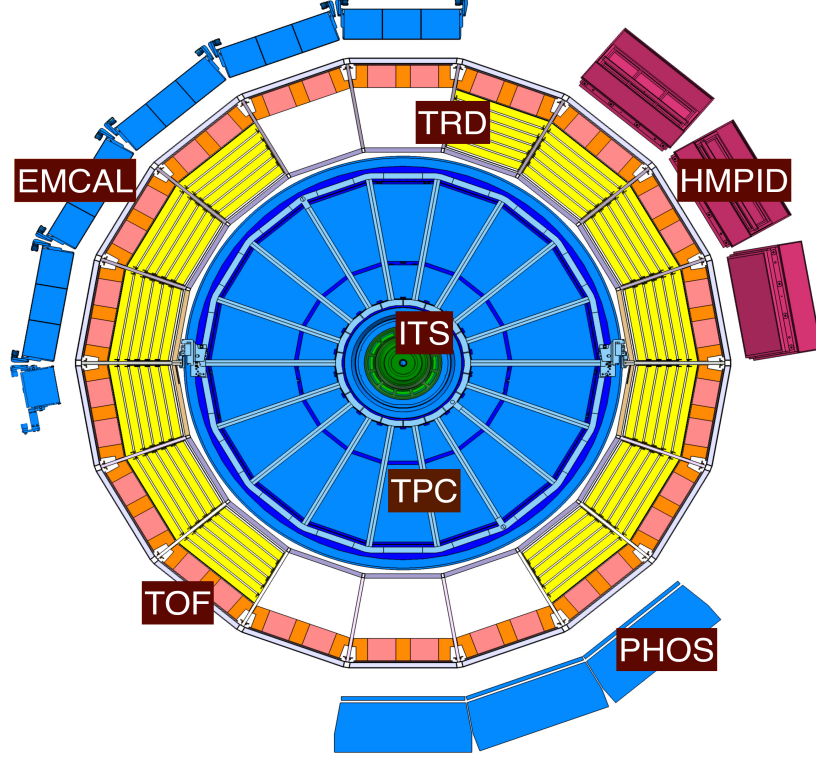


Figure 2.3: (Color online) Lateral view of The ALICE central barrel detector at the end of 2012 [29].

Besides tracking, ITS and TPC helps in charged-particle identification via measurement of the specific ionization energy loss dE/dx . The TRD detector consists of six layers of Xe–CO₂-filled MWPCs which has a fiber/foam radiator in front of each chamber. The main work of TRD is to track charged particles and also identify electron via transition radiation and dE/dx . The TOF detector is based on Multigap Resistive Plate Chamber (MRPC) technology is used for particle identification at intermediate momenta through the time of flight of particles. Finally, outside TOF the cylindrical volume is shared by two electromagnetic calorimeters with thicknesses of $\sim 20 X_0$ (radiation lengths) and $\sim 1 \lambda_{\text{int}}$ (nuclear interaction length), the high-resolution PHOTon Spectrometer (PHOS) and the large-acceptance ElectroMagnetic CALorimeter (EMCal), along with the

ring-imaging Cherenkov detector High Momentum Particle Identification Detector (HMPID), which has a liquid C_6F_{14} radiator and a CsI photo-cathode for charged-hadron identification at intermediate momenta.

2.2.2 Forward detectors

The ALICE forward detectors are placed in the forward rapidity regions and covered by the preshower Photon Multiplicity Detector (PMD) and the silicon Forward Multiplicity Detector (FMD). These two detectors are dedicated to the measurement of photons and charged particles around $|\eta| \approx 3$, respectively. The T0 is a quartz Cherenkov detector delivers the time and the longitudinal position of the interaction whereas the V0 is a plastic scintillator detector with rapidity coverage at $-3.7 < \eta < -1.7$ and $2.8 < \eta < 5.1$, mainly used for triggering and centrality determination through the charged particle distribution. The Zero Degree Calorimeter (ZDC) consists of two tungsten-quartz neutron (ZN) and two brass-quartz proton (ZP) calorimeters, located symmetrically on both sides at about 116 m from of the interaction point, which used to count spectator nucleons and also helps in centrality determination.

2.2.3 MUON spectrometer

The MUON spectrometer is used to measure quarkonium and light vector meson production in a region of $-4.0 < y < -2.5$. It is also used to study high- p_T muons predominantly coming from the decay of charm and beauty. The MUON spectrometer is made of a hadron absorber of $\sim 10 \lambda_{\text{int}}$, a dipole magnet of 3 Tm, and five tracking stations with two pad chambers each (Muon Chambers, MCH), helps in measuring quarkonium and light vector meson production in a region of $-4.0 < y < -2.5$. Single-muon and muon-pair triggers with an adjustable transverse-momentum threshold are given by two further stations (Muon Trigger, MTR) placed behind an additional $7\lambda_{\text{int}}$ absorber.

The K^{*0} and ϕ mesons are reconstructed in the mid-rapidity with the advantage of unique tracking and particle identification of ALICE central barrel detectors. The relevant detectors for this analysis are the ITS, the TPC and TOF. In addition to this V0AND is used for triggering minimum bias events.

2.3 Particle Tracking and Identification via Central Barrel Detector

2.3.1 The Inner Tracking System (ITS)

The main objectives of ITS are to localize the primary vertices with a resolution better than $100\ \mu\text{m}$ and also helps in reconstructing the secondary vertices from decays of hyperons, B, and D mesons. In addition ITS helps in tracking and identification of charged particles with transverse momentum p_T below $200\ \text{MeV}/c$ or particles traversing the dead zones of the TPC, which improves the momentum and angle resolution for particles reconstructed by the TPC. Fig. 2.4 shows the schematic diagram of ITS, which is composed of 6 cylindrical layers of silicon detectors located at radii between 3.9 and 43 cm and surrounds the beam pipe which is a $800\ \mu\text{m}$ thick beryllium cylinder with 2.9 cm radius. It has a rapidity coverage of $|\eta| < 0.9$ for all vertices located within the length of the interaction diamond which is $\pm 5.3\ \text{cm}$ along the beam direction.

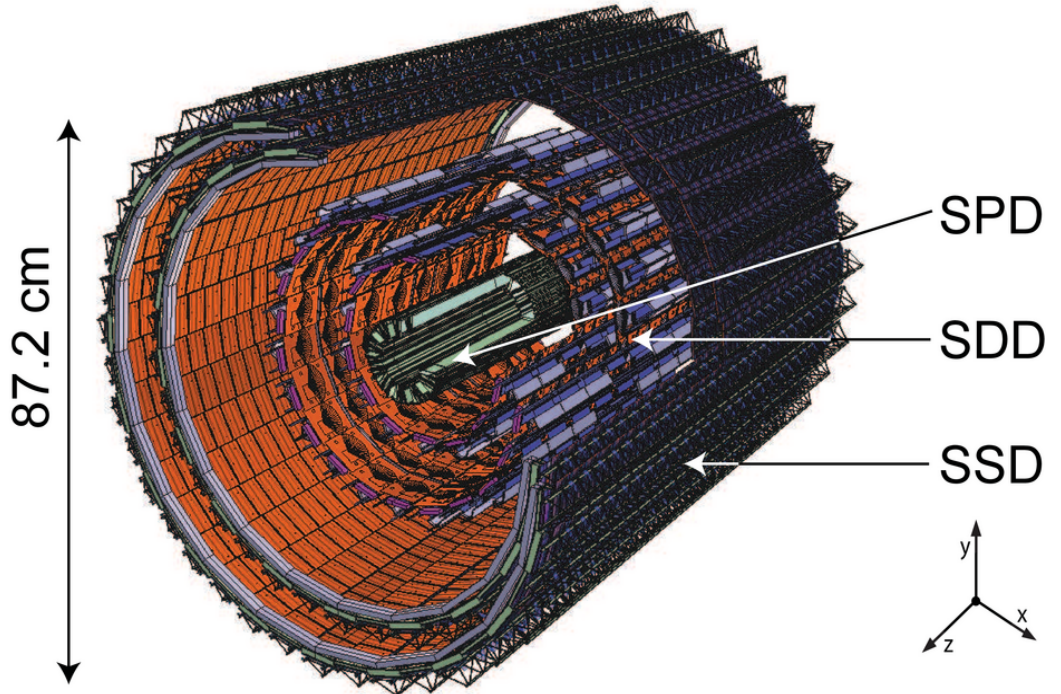


Figure 2.4: (Color online) ALICE Inner Tracking System [30].

The first two innermost ITS layers constitute the Silicon Pixel Detectors

(SPD) with average radii 3.9 and 7.6 cm. SPD consists of 9.8×10^6 pixels with dimension $50 (r\varphi) \times 425 (z) \mu\text{m}^2$ which has a spatial resolution of $12 (r\varphi) \times 100 (z) \mu\text{m}$. The high granularity of SPD helps to determine the primary vertex with a resolution of $100 \mu\text{m}$ and also helps in minimum bias triggering.

The two intermediate layers with average radii 15 and 23.9 cm are made up of Silicon Drift Detectors (SDD). The third and fourth layers have a spatial resolution of $35 \mu\text{m}$ in the $r\varphi$ direction and $25 \mu\text{m}$ in the z direction. It has a sensitive area of 70.2 mm^2 in $r\varphi$ and 75.3 mm^2 in z direction which splits into two drift regions, where electrons move in opposite direction under a drift field of 500 V/cm . The SDDs modules are mounted on linear structures known as ladders. The layer 3 has 14 ladders per layer with 6 modules each, whereas layer 4 has 14 ladders with each of them having 8 modules. The position of particle in the z -direction from the centroid is reconstructed from the charge collected along anode, whereas the drift time with respect to the trigger time gives r coordinate along the drift.

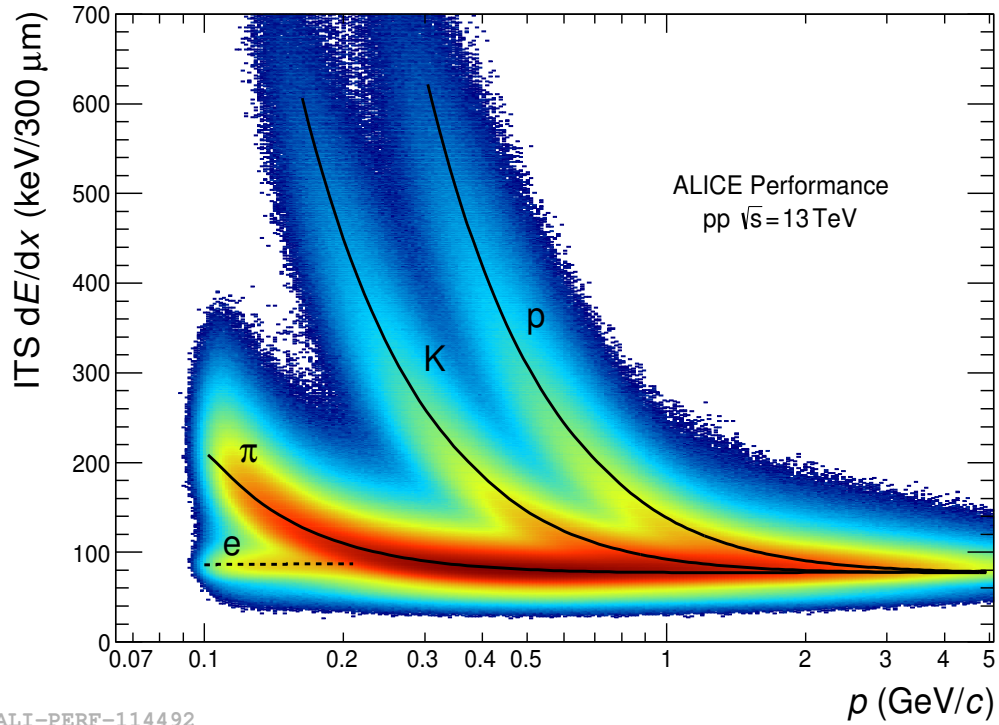


Figure 2.5: (Color online) Average dE/dx of charged particles vs their momentum for ITS pure standalone tracks measured in pp collisions at 13 TeV.

The two outermost ITS layers are made up of Silicon Strip Detectors (SSD). Each SSD module is composed of double sided strip sensor with an intrinsic spatial resolution of $20\text{ }r\varphi \times 830\text{ } (z)\text{ }\mu\text{m}^2$. The sensors have an active area of $73\text{ } (r) \times 40\text{ } (z)\text{ }mm^2$ with a thickness of $300\text{ }\mu\text{m}$ and contains 768 strips on each side with a pitch of $95\text{ }\mu\text{m}$. The innermost SSD layer consists of 34 ladders, each of them having 22 modules along the beam direction whereas other SSD layers are made up of 38 ladders, each of them having 25 modules. The first four outer layers have analogous readouts which has a large dynamic range that allows the particle identification via energy loss (dE/dx) in the non-relativistic ($1/\beta^2$) region for low momentum particles down to $p_T = 100\text{ MeV}/c$. Fig. 2.5 shows the average energy loss (dE/dx) of charged particles vs their momentum, both measured by the ITS alone (ITS pure standalone track) in pp collisions at 13 TeV LHC15f pass2 (data taken in 2015) period (ITS pure standalone reconstruction). The lines are a parametrization of the detector response based on a hybrid parametrization with a polynomial function at low p/m and a “PHOBOS” Bethe-Bloch formula. The plots shows the capability of ITS for particle identification via energy loss (dE/dx), where electrons can be identified from $100\text{ MeV}/c$ to $200\text{ MeV}/c$. Pions are separated from kaons up to $0.6\text{ GeV}/c$ and similarly kaons from protons up to $0.8\text{ GeV}/c$.

The momentum and angular resolution of tracks reconstructed by TPC can be improved by ITS and also allows to reconstruct the tracks traversing the dead TPC region or very low p_T tracks that do not reach the TPC. ALICE take the advantage of ITS capabilities with a dedicated tracking algorithms which helps to recover very low momentum particles. The main difference between ITS standalone tracking and global tracking is the starting point. The ITS standalone tracking starts from the primary vertex and it reconstruct the tracks by matching the clusters coming from the different ITS layers through selected angular windows.

2.3.2 The Time Projection Chamber (TPC)

The TPC is a large 3-dimensional tracking device and the heart of ALICE central barrel tracking. The TPC has a pseudorapidity coverage of $|\eta| < 0.9$ with

full azimuthal coverage (with exception of the dead zone) is optimised to provide charged particle momentum measurements with good two-track separation, particle identification and vertex determination together with the other central barrel detectors. The TPC is optimized to provide better tracking performance in high multiplicity environments, hence the material budget is kept as low as possible in order to have very low multiple scattering and secondary particle production. The TPC provides a large cylindrical volume with an inner radius of 80 *cm*, determined by the maximum acceptable track density and with a outer radius of 250, *cm* determined by the minimum track length at which the resolution on dE/dx is lower than 10%. The TPC has an active volume of 92 m^3 and is filled with neon (Ne) and carbon-dioxide (CO₂) in which the primary electrons are transported over a distance up to 2.5 *m* on either side of the central electrode to the end plates. The TPC is divided into two drift regions by a central cathode and the end plates consist of multi-wire proportional chambers (MWPC). The MWPC with cathode pad readout are mounted in to 18 trapezoidal sectors at each end plate. The ionised electrons by a charged particle drift towards the end plates with a drift velocity of 2.7×10^8 *cm/s* gives a maximum drift time of 92 μs , resulting a limit on the event rate sustainability of the TPC. The pile-up effects become more relevant at high interaction rate and tracks from pile-up events can be rejected using cuts on primary vertex position and arrival time, and these factors force ALICE to run at lower instantaneous luminosity with respect to other LHC experiments. In addition to this TPC has high busy time that slows down the read out frequency. The TPC has a radial thickness of 3.5 % of the radiation length (X_0) at central rapidity and grows to 40 % towards the edges.

The TPC has a low material budget and rather low magnetic field (0.5 T) are the requirements to optimize both tracking and particle identification down to very low momentum ($p \geq 100$ MeV/*c*) with a good momentum resolution and efficiency > 80 % for $p_T > 100$ MeV/*c* as reported in Fig. 2.6. By measuring the deflection in the magnetic field, ITS along with TPC can able to determine the momentum of the charged particles with a resolution better than 1 % at low p_T and better than 20 % for $p_T \simeq 100$ MeV/*c*.

Inside the TPC, a charged particle traversing the TPC induces signals on

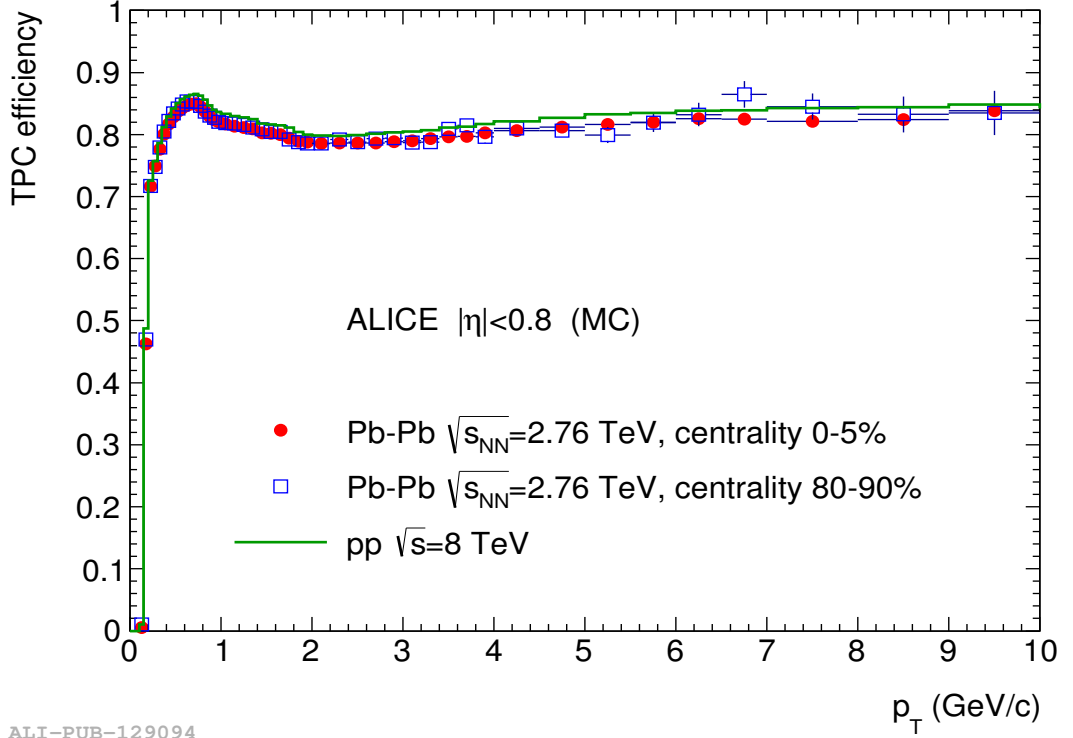


Figure 2.6: (Color online) TPC track finding efficiency for primary particles in MC for pp and Pb-Pb collisions.

a maximum of 159 tangential pad-rows, each corresponding to one cluster used in track reconstruction, which corresponds to the total number of pad rows in a given TPC sector. The particle identification (PID) through TPC is done by simultaneously measuring the particle energy loss in the TPC readout pads, charge and momentum of each particle while traversing through the TPC gas. The specific energy loss in TPC is described by the Bethe-Bloch formula and parametrised by the following function,

$$-\left\langle \frac{dE}{dx} \right\rangle = k_1 z^2 \frac{Z}{A} \frac{1}{\beta^2} \left[\frac{1}{2} \ln \left(k_2 m_e c^2 \beta^2 \gamma^2 \right) - \beta^2 + k_3 \right] \quad (2.5)$$

here $\beta\gamma = p/Mc$ and γ is the Lorentz factor, β is the particle velocity, Z is the atomic number of the ionised gas, A is the mass of the ionised gas, m_e is the electron mass, z is electric charge of the ionising particle, M is the ionising particle mass, p is momentum of the ionising particle, and k_{1-3} are the fitting parameters depend on the ionised medium. Fig. 2.7 shows the specific energy loss (dE/dx) as function of particle momentum in the TPC in pp collisions at

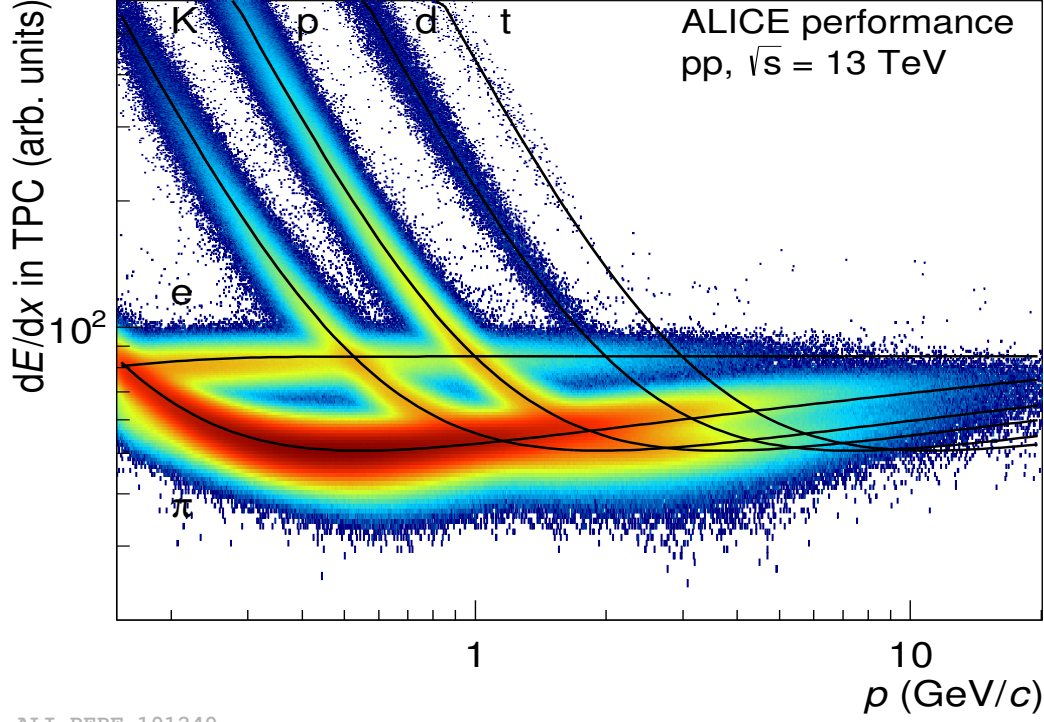


Figure 2.7: (Color online) TPC track finding efficiency for primary particles in MC for pp and Pb-Pb collisions.

$\sqrt{s} = 13$ TeV and exhibit a clear separation between different particle species (The solid line shows the expected value of specific energy loss ($\langle \frac{dE}{dx} \rangle$) for e, π , K, p and deuteron).

In non-relativistic region, the specific energy loss decrease with momentum due to the increasing value of β until $\beta \simeq 0.96$ and reaches a minimum value. Here the minimum value of the dE/dx is almost same for the charged particles having same charge and the particle at this point is known as minimum ionising. Due to the well separated bands of the dE/dx curve, particles at lower momentum ($p \leq 1$ GeV/c) can be identified on track-by-track basis. Similarly, for higher momentum the dE/dx bands for different particles start to overlap and can be separated by a statistical basis via multi-Gaussian fit as shown in Fig. 2.8. With long tracks (≥ 130 samples) and with the truncated-mean method the resulting dE/dx peak shape is Gaussian. In the relativistic rise region the dE/dx shows a constant separation between different bands in a wide range of momentum.

In the TPC different particles can be identified by applying a cut on the N_σ

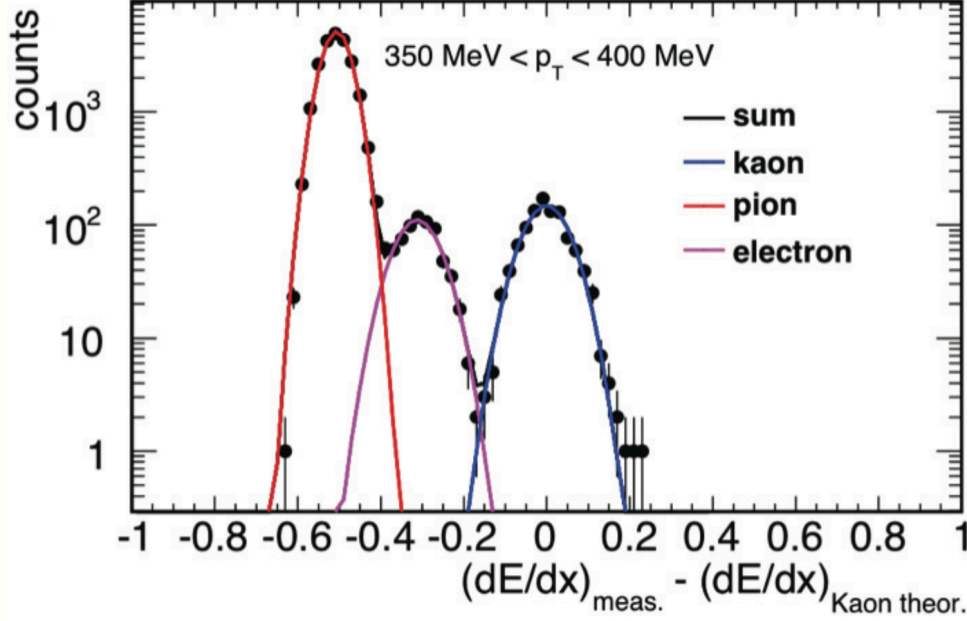


Figure 2.8: (Color online) The ionization signal in a momentum window of 50 MeV/ c width in the $1/\beta^2$ region of the Bethe-Bloch curve.

value, which is the difference between the specific energy loss in the experiment to the theoretical value calculated from Bethe-Bloch formula scaled with the TPC resolution and given as

$$N_\sigma = \frac{(dE/dx)_{measured} - (dE/dx)_{Bethe-Bloch}}{\sigma_{TPC}}. \quad (2.6)$$

2.3.3 The Time of Flight (TOF)

The Time-Of-Flight (TOF) detector is based on Multi-gap Resistive Plate Chambers (MRPCs) technology and has 1593 glass Multi-gap Resistive Plate Chamber (MRPC) detectors, each with a sensitive area of $7.4 \times 120 \text{ cm}^2$. The TOF covers the pseudo-rapidity interval $|\eta| < 0.9$ with full azimuthal angle acceptance, with a total active area of 141 m^2 . Each MRPC is composed of 96 readout pads of area $2.5 \times 3.5 \text{ cm}^2$ and in total consists of 152928 readout channels. This highly segmented structure helps to have a low occupancy and a good performance also at high multiplicities.

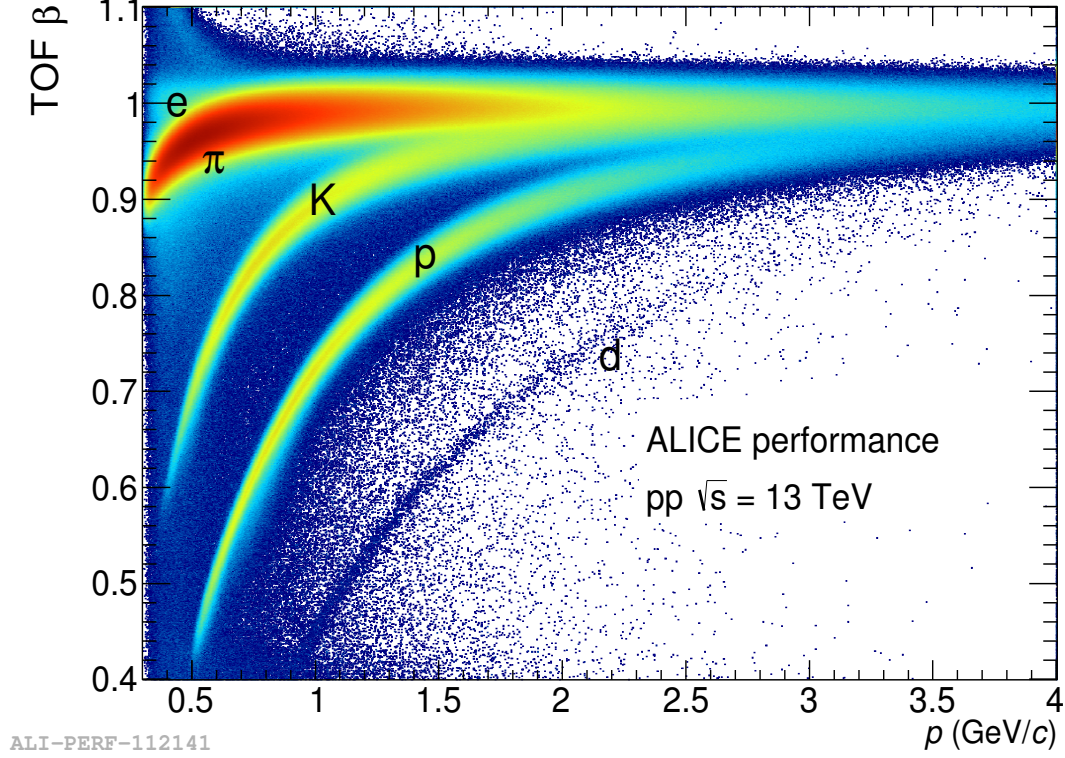


Figure 2.9: (Color online) TOF beta vs Momentum plot in pp collisions at $\sqrt{s} = 13$ TeV.

The TOF detector identify particles by combining the measurement of particle time-of-flight with the momentum information provided by the TPC detector. A time resolution of 100 ps is needed to separate pions and kaons up to 2.5 GeV/ c and protons up to 4 GeV/ c . For the time-of-flight measurement the start time is provided by the T0 detector, which consists of two arrays of Cherenkov counters T0C ($-3.28 < \eta < -2.97$) and T0A ($4.61 < \eta < 4.92$), located at opposite sides of the interaction point (IP). The arrival time can also be estimated from the arrival time of TOF with the conditions of three particles should reach the TOF detector. This is more important for the cases, where T0 detector is absent. Events where these two methods are absent, an average TOF start time for the run is used.

Fig 2.9 presents the ALICE TOF performance in pp at $\sqrt{s} = 13$ TeV, and here the particle velocity is shown as a function of momentum measured by TPC. Particles are identified in the TOF detector by comparing the measured time of flight to the expected time for a given particle species and this is also exhibits in

Fig. 2.9. Tracks that are incorrectly matched with the TOF contribute towards background and it is due to the higher track density.

Chapter 3

Measurement of K^{*0} and ϕ Mesons in pp Collisions

The study of resonances helps us to understand the particle production mechanisms. Resonance productions at LHC energies involve soft and hard-scattering processes. The hard scatterings are perturbative in nature and responsible for production of high- p_T (≥ 2 GeV/ c) particles, whereas the bulk of the particles are produced due to soft scattering processes which are non-perturbative in nature. The contributions from high- p_T particles originate from fragmentation of jets and can be calculated folding the perturbative QCD calculations for jets with universal fragmentation functions determined from experimental data. The production of low- p_T particles require phenomenological models due to non-perturbative regime which can not be estimated from the first principles of QCD. In this chapter, we discuss about $K^{*0}(892)$ and $\phi(1020)$ which have lifetimes 4.16 ± 0.05 fm/ c and 46.3 ± 0.4 fm/ c , respectively [32]. The $\phi(1020)$ is one of the lightest vector mesons consisting of $(s\bar{s})$, helps in study of the strangeness production. The production of $s\bar{s}$ was found to be significantly suppressed while comparing to $u\bar{u}$ and $d\bar{d}$ pairs in pp collisions [34, 35]. The $K^{*0}(892)$ is also a vector meson having a similar mass to $\phi(1020)$, but differs in strangeness content by one unit which also helps to probe the strangeness production. The measurement of particle production in minimum bias pp collisions provide an input to tune the perturbative-Quantum Chromodynamics (pQCD) inspired Monte Carlo (MC) event generators such as PYTHIA [36] and EPOS [37]. Furthermore, the measurements in minimum bias

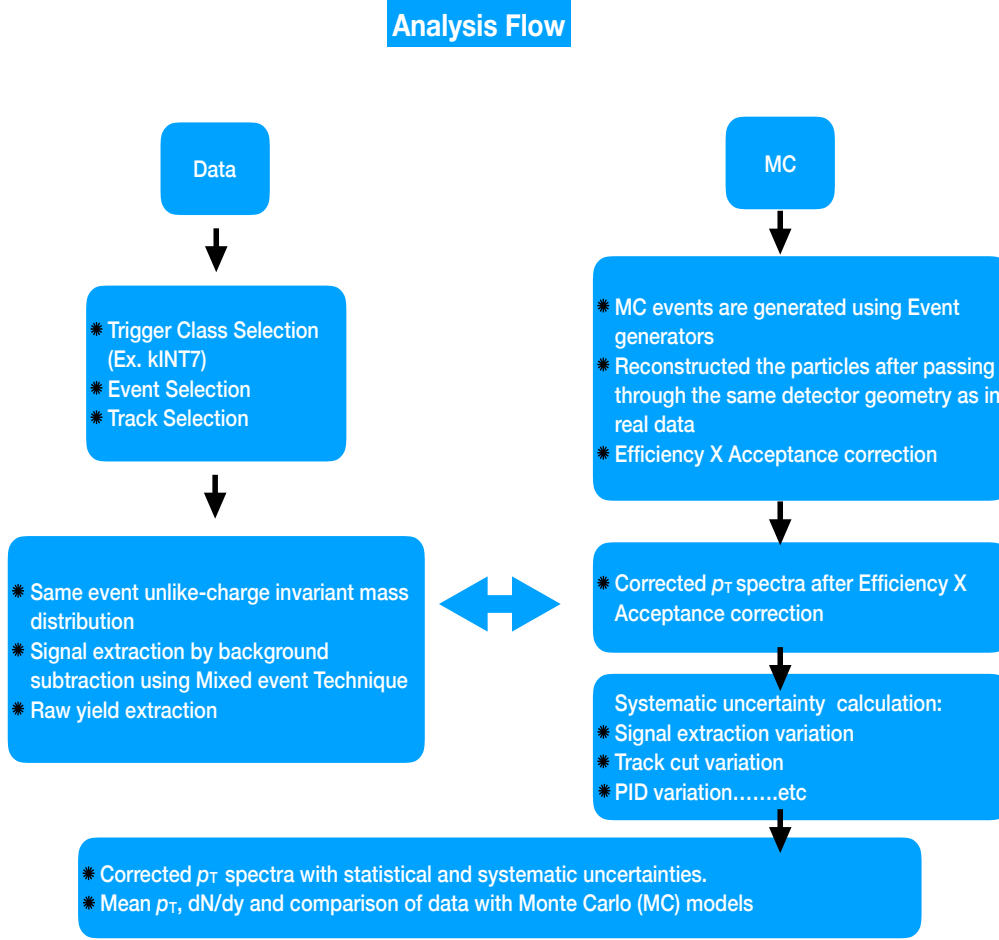


Figure 3.1: (Color online) Analysis flow chart.

pp collisions at $\sqrt{s} = 8$ TeV reported in this analysis serve as reference data to study nuclear effects in proton-lead (p-Pb) and lead-lead (Pb-Pb) collisions. The flow chart for a quick overview of the sequential movement of the analysis is shown in Fig. 3.1.

3.1 Data Set and Event Selection

The data from proton-proton collisions at $\sqrt{s} = 8$ TeV (recorded in 2012) are analyzed by taking AOD data from LHC12a, LHC12b, LHC12c, LHC12d, LHC12h and LHC12i pass 2. After the event selection the accepted events are around 46 M events for K^{*0} and around 51 M events for ϕ , selected using the kINT7 trigger which will be discussed later in this section.

The following runs are selected from the Run Condition Table (RCT) with global quality good (flag 1, 25) and the detectors required for this selection are SSD, SPD, SDD, V0, TPC, TOF and T0.

* **LHC12a:** 176701, 176704, 176715, 176730, 176749, 176752, 176753, 176849, 176854, 176859, 176924, 176926, 176929, 177011, 177157, 177160, 177167, 177173, 177180, 177182.

* **LHC12b:** 177864, 177860, 177858, 177804, 177799, 177798, 177681, 177680, 177679, 177624, 177620, 177612, 177597, 177592, 177580.

* **LHC12c:** 182744, 182741, 182740, 182729, 182728, 182725, 182724, 182692, 182691, 182687, 182684, 182635, 182513, 182509, 182325, 182324, 182323, 182322, 182302, 182300, 182299, 182297, 182295, 182289, 182207, 182111, 182110, 182106, 182023, 182022, 182018, 182017, 180720, 180719, 180717, 180716, 180569, 180567, 180561, 180517, 180515, 180510, 180501, 180500, 180230, 180201, 180200, 180199, 180133, 180132, 180131, 180130, 180129, 180127, 180044, 180042, 180000, 179920, 179919, 179918, 179917, 179916, 179859, 179858, 179803, 179802, 179796, 179639, 179638, 179621, 179618.

* **LHC12d:** 186320, 186319, 186208, 186205, 186167, 186165, 186164, 186163, 185784, 185778, 185776, 185775, 185768, 185765, 185764, 185738, 185687, 185589, 185588, 185583, 185582, 185581, 185580, 185578, 185575, 185574, 185465, 185461, 185375, 185371, 185363, 185362, 185361, 185360, 185359, 185356, 185351, 185350, 185349, 185303, 185302, 185300, 185299, 185296, 185293, 185292, 185291, 185289, 185288, 185284, 185282, 185221, 185217, 185208, 185206, 185203, 185198, 185196, 185189, 185164, 185160, 185157, 185134, 185132, 185127, 185126, 185116, 185031, 185029, 184928, 184786, 184784, 184687, 184682, 184678, 184673, 184215, 184209, 184208, 184188, 184138, 184137, 184135, 184132, 184131, 184127, 184126, 183916, 183913.

* **LHC12h:** 192732, 192731, 192729, 192510, 192505, 192499, 192492, 192468, 192461, 192453, 192417, 192415, 192349, 192348, 192347, 192246, 192205, 192202, 192201, 192200, 192199, 192194, 192177, 192174, 192140, 192128, 192095, 192075, 192073, 192072, 192004, 191244, 191234, 191231, 191229,

191227, 190984, 190983, 190981, 190979, 190975, 190974, 190970, 190968,
 190904, 190903, 190895, 190425, 190421, 190418, 190417, 190416, 190393,
 190392, 190390, 190389, 190386, 190342, 190341, 190340, 190338, 190337,
 190216, 190215, 190214, 190212, 190210, 189698, 189697, 189696, 189621,
 189616, 189612, 189610, 189605, 189603, 189577.

* **LHC12i:** 193151, 193094, 193051, 193049, 193047, 193014, 193011, 193010,
 193008, 193007, 193005, 193004, 192824, 192822, 192820, 192779, 192778,
 192775, 192772.

For this analysis, we select good events which require a coincidence signal in V0-A and V0-C plastic scintillators and will be discussed later in this section. Events having the primary vertex coordinate along the beam axis within 10 cm from the interaction point are selected for this analysis. A pileup rejection on the events are applied when more than one vertex is found in the Silicon Pixel Detector (SPD), the two innermost layers of the Inner Tracking System (ITS) . Figure 3.2 shows the event statistics after implementing all the above event-selection cuts. The fourth bin shows the total number of accepted events for this analysis, which is around 46 M events for K^{*0} out of 49 M triggered events (CINT1B). Let us discuss about the minimum bias trigger (kINT7) used for this analysis.

Minimum Bias Trigger (kINT7)

The kINT7 trigger uses a stricter condition using the coincidence signal between both the VZERO arrays during the 2011 and 2012 pp data taking periods, with the increasing LHC luminosity and beam background. Also it helps to reject the interaction of the beam with the residual gas in the beam pipe. The VZERO detector is based on the principle that the arrival time of produced particles to V0A and V0C will be different in pp and beam-gas interactions.

Figure 3.3 shows the schematic diagram of arrival time of particles to the V0A and V0C counters for three different conditions. Let t_0 be the time when the two proton bunches cross each other in the nominal interaction point and this time passed on from the LHC clock to the VZERO trigger.

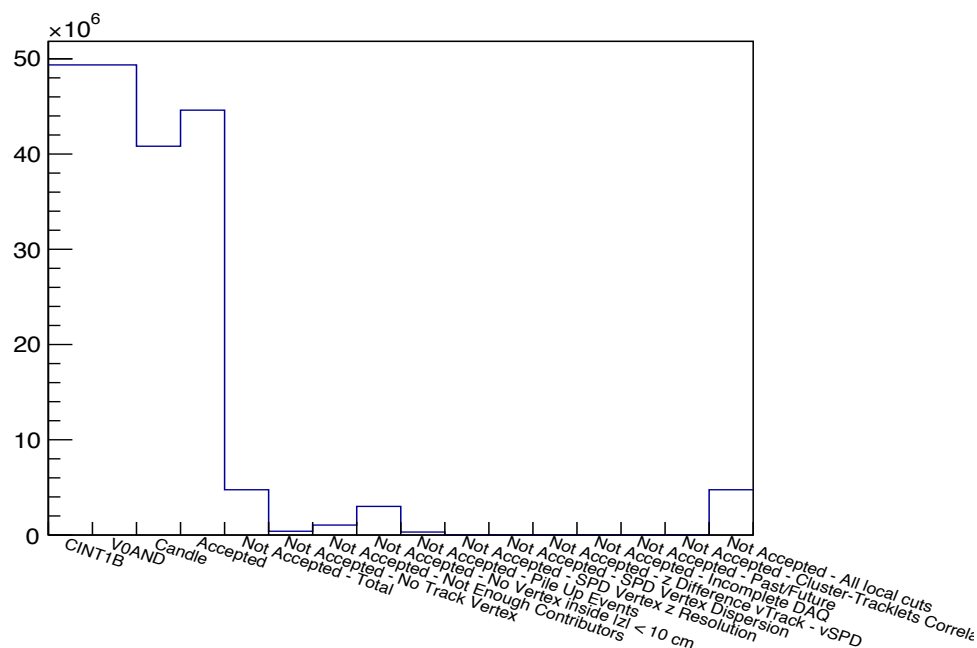


Figure 3.2: Event statistics after implementing event-selection cuts. The fourth bin shows the total number of accepted events for this analysis [38–40].

- (a) The arrival time of particles coming from the pp collisions (Figure 3.3 (a)) to V0A (V0C) array is approximately 11.3 ns (3.0 ns) after t_0 .
- (b) Particles originating from the beam-background behind the V0C will arrive at V0C around 3.0 ns before t_0 and 11.3 ns after t_0 at V0A.
- (c) Similarly particles originating from the beam-background behind the V0A will arrive at V0A around 13.3 ns before t_0 and 3.0 ns after t_0 at V0A.

In these cases the arrival time to the VZERO arrays relative to t_0 are not exact rather has a spatial distribution and is due to (i) the collision does not happen exactly near the nominal interaction point but a region around it, (ii) the collision happens not exactly at t_0 but in a time region around it and other factors. The correlation of sum and differences of arrival times of particles to both the VZERO arrays results in three different types of event classes as shown in Fig 3.4.

- (1) Collisions at the nominal interaction points with coordinates at (8.3 ns, 14.3 ns).

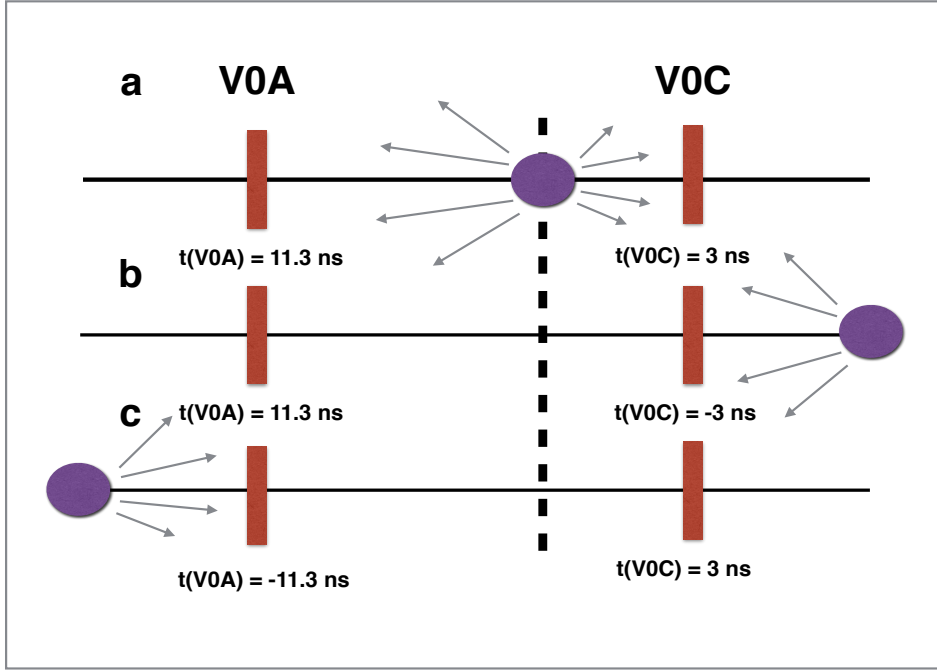


Figure 3.3: (Color online) (a) A schematic diagram of arrival time of particles to the V0A and V0C counters in pp collisions, (b) for beam-background interactions from C-side and (c) for beam-background interactions from A-side. Here $t(V0A)$ and $t(V0C)$ are the arrival times of particles at V0A and V0C, respectively.

(2) Beam-background collisions from C side at (14.3 ns, 8.3 ns).

(3) Beam-background collisions from A side at (-14.3 ns, -8.3 ns).

Fig 3.4 shows a clear separation of background events in two peaks in the time sum-difference plane of signals from the main collision peak. V0 time windows used to set the trigger conditions on collision events help for the selection of good events.

Now let's discuss about the track selection and particle identification criteria used in the present analysis.

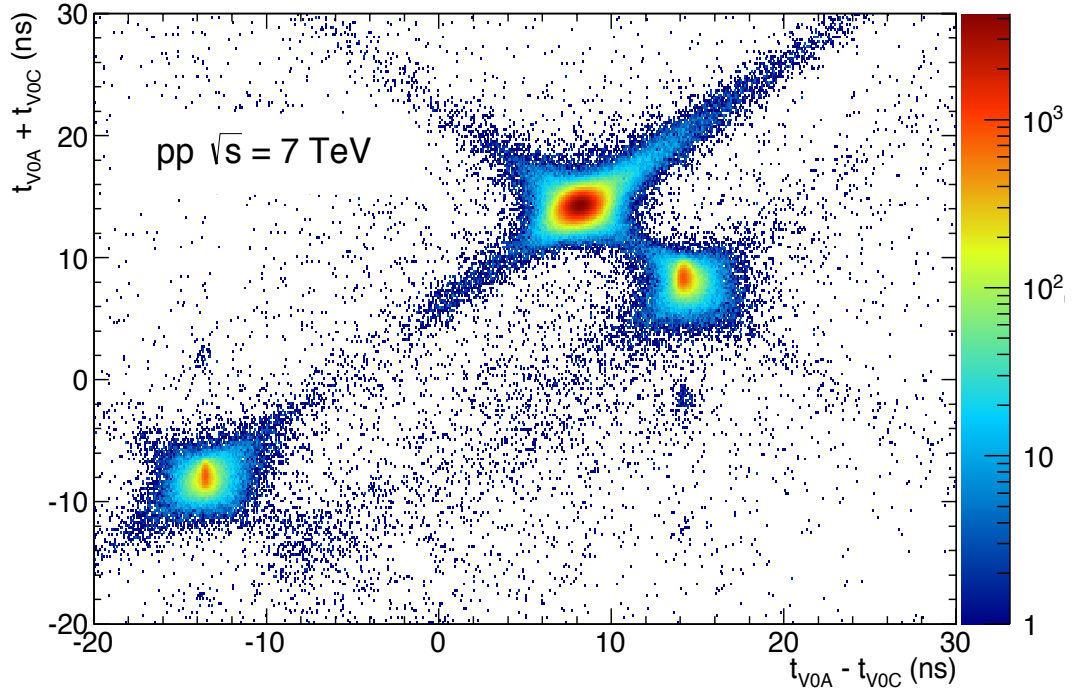


Figure 3.4: (Color online) Correlation of sum and differences of arrival times of the particles to the both the VZERO arrays results in three different types of event classes.

3.2 Track Selection and Particle Identification

The K^{*0} and ϕ mesons are identified by reconstructing with their hadronic decays. The K^{*0} is reconstructed from $\pi^\pm$ and $K^\mp$ pairs with 66 % branching ratio and similarly the ϕ is reconstructed from Kaon pairs ($\phi \rightarrow K^+ K^-$) having a branching ratio of 0.489 ± 0.005 [32]. A particle traversing the TPC induces signals on a maximum of 159 tangential pad-rows, each corresponding to one cluster used in track reconstruction, which corresponds to the total number of pad rows in a given TPC sector. A track is considered as good if it induces signals with at least 70 clusters and this selection is applied for pion and kaon tracks from the decays of K^{*0} and ϕ . For a good track the χ^2 per track point ($\chi^2/N_{\text{clusters}}$) of the track fit in the TPC must be less than 4. Tracks must be associated with at least one cluster in the SPD and the χ^2 per cluster in the ITS, which is limited to 36. Tracks are selected within $|\eta| < 0.8$ to ensure a uniform acceptance by avoiding the edges of the TPC. Cuts on the distance of closest approach to the

primary vertex in the transverse plane (DCA_{xy}) and longitudinal direction (DCA_z) are applied in order to reduce contamination from beam-background events and secondary particles coming from weak decays. The value of DCA_{xy} is required to be less than 7 times its resolution: $DCA_{xy}(p_T) < (0.0105 + 0.035p_T^{-1.1})$ cm (p_T in GeV/ c) and DCA_z is required to be less than 2 cm. To improve the global resolution, the p_T of each track is restricted to be greater than 0.15 GeV/ c . The track cut criteria are summarised below

1. minimum number of crossed rows in TPC: $N_{\text{cr,TPC}} \geq 70$
2. minimum ratio of number of crossed rows to number of findable clusters in TPC: $R_{\text{TPC}} \geq 0.8$
3. maximum χ^2 per cluster in TPC: $\chi_{\text{TPC}}^2 < 4$
4. reject kink daughters
5. require ITS refits
6. require TPC refits
7. minimum number of clusters in SPD: 1
(`AliESDtrackCuts::SetClusterRequirementITS(kSPD, kAny)`)
8. DCA_{xy} p_T formula: $DCA_{xy} < 0.0105 + 0.035p_T^{-1.1}$ (a 7- σ cut)
9. maximum χ^2 in TPC-Constrained Global: $\chi_{\text{CGI}}^2 < 36$
10. $DCA_z < 2$ cm
11. `AliESDtrackCuts::SetDCAToVertex2D(kFALSE)`
12. `AliESDtrackCuts::SetRequireSigmaToVertex(kFALSE)`
13. maximum χ^2 per cluster in ITS: $\chi_{\text{ITS}}^2 < 36$
14. $p_{T,\text{track}} > 0.15$ GeV/ c
15. $|\eta_{\text{track}}| < 0.8$

Charged particle identification has been performed using the TPC and TOF detectors. In the TPC, particles are identified by the specific ionization energy loss (dE/dx) in the TPC gas, whereas for the TOF it is done by measuring the time of flight. In the TPC particle species are selected using a cut on the difference of the mean values of the specific energy loss to the expected dE/dx values divided by the resolution σ_{TPC} which also reduces the contamination from other charged particles. As described below, this is optimized for each analysis and depends on the signal-to-background ratio and on the transverse momentum. Particles are identified in the TOF by comparing the measured time of flight to the expected time for a given particle species. The cut is expressed in units of the estimated resolution σ_{TOF} , which has a mean value of 160 ps [41]. Thanks to the time resolution of TOF which allows pions and kaons to be unambiguously identified up to momentum $p \approx 1.5 \text{ GeV}/c$ and also removes contamination from electrons. The two mesons (pions and kaons) can be distinguished from (anti)protons up to $p \approx 2.5 \text{ GeV}/c$.

Three TPC-PID selection criteria are used for K^{*0} and ϕ reconstruction depending on the momentum of the daughter particles. Both pions and kaons are selected using a cut of $|N\sigma_{\text{TPC}}| < 2.0$ for $p(K^\pm, \pi^\pm) > 0.4 \text{ GeV}/c$. For $p(K^\pm, \pi^\pm) < 0.3 \text{ GeV}/c$, a cut of $|N\sigma_{\text{TPC}}| < 6.0$, while a cut of $|N\sigma_{\text{TPC}}| < 4.0$ for $0.3 < p(K^\pm, \pi^\pm) < 0.4 \text{ GeV}/c$ is applied. Also, a TOF 3σ veto cut is applied for K^{*0} . “TOF veto” means that the TOF 3σ cut is applied only for cases where the track matches a hit in the TOF. For tracks without matches, no TOF cut is applied. Fig. 3.5 and 3.6 show the two different TPC PID cut schemes i.e, (left panel) without any PID cut, whereas (right panel) with the tpc2tof3 scheme. Similarly, Fig. 3.7 and 3.8 show without the TOF PID and with a TOF 3σ cut for both pions and kaons.

3.3 Raw Yield Extraction

The K^{*0} (ϕ) meson is reconstructed through its dominant hadronic decay channel $K^{*0} \rightarrow \pi^\pm K^\mp$ ($\phi \rightarrow K^+ K^-$). The invariant mass is constructed by taking unlike-sign pairs of K and π (K) candidates for K^{*0} (ϕ) in the same event. The rapidity

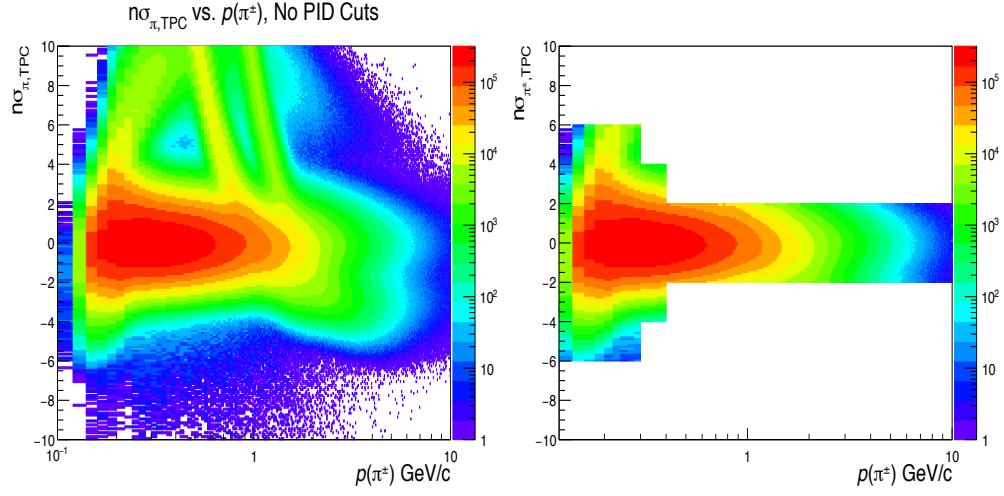


Figure 3.5: (Color online) $n\sigma_{\pi,TPC}$ vs. p for two different PID cut schemes. **Left Panel:** no PID cuts applied. **Right Panel:** default PID cuts (the tpc2tof v3 scheme) applied [38–40].

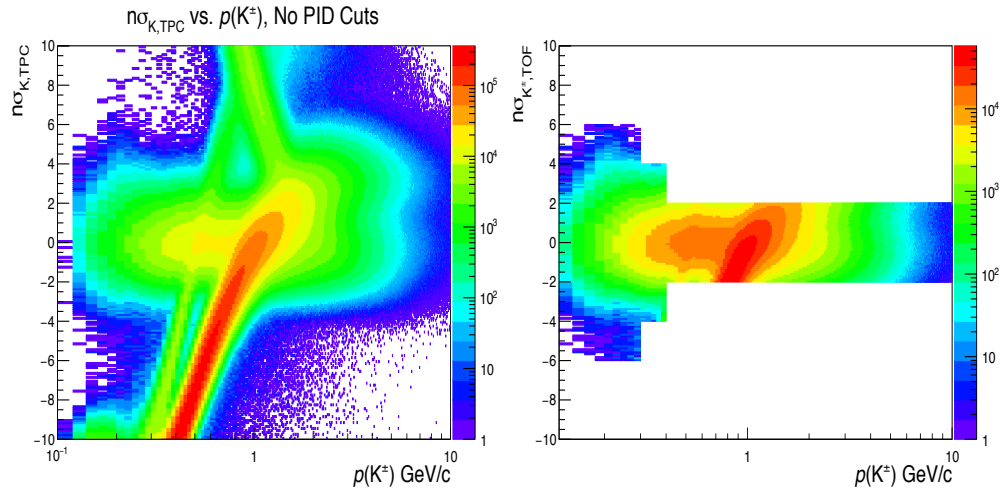


Figure 3.6: (Color online) $n\sigma_{K,TPC}$ vs. p for two different PID cut schemes. **Left Panel:** no PID cuts applied. **Right Panel:** default PID cuts (the tpc2tof v3 scheme) applied [38–40].

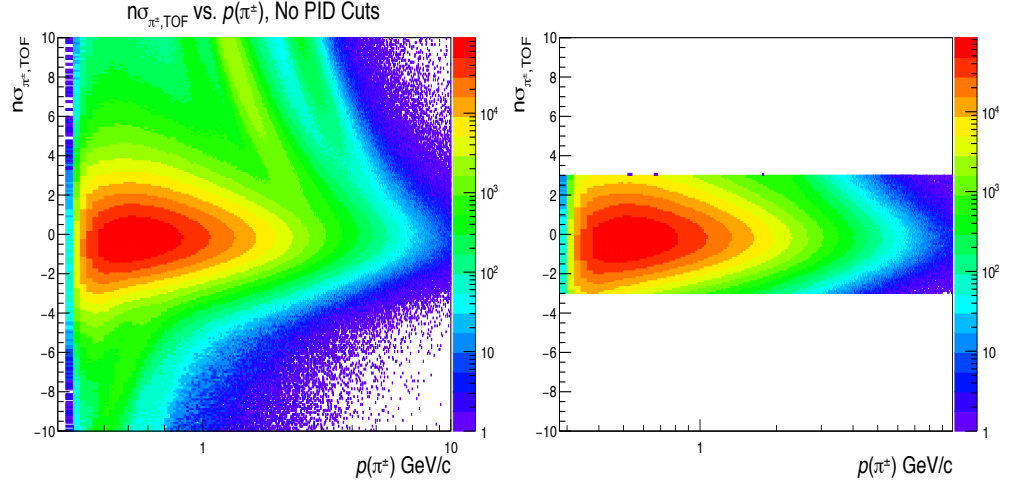


Figure 3.7: (Color online) $n\sigma_{\pi, TOF}$ vs. p for two different PID cut schemes. **Left Panel:** no PID cuts applied. **Right Panel:** a $|n\sigma_{\pi, TOF}| < 3$ cut applied [38–40].

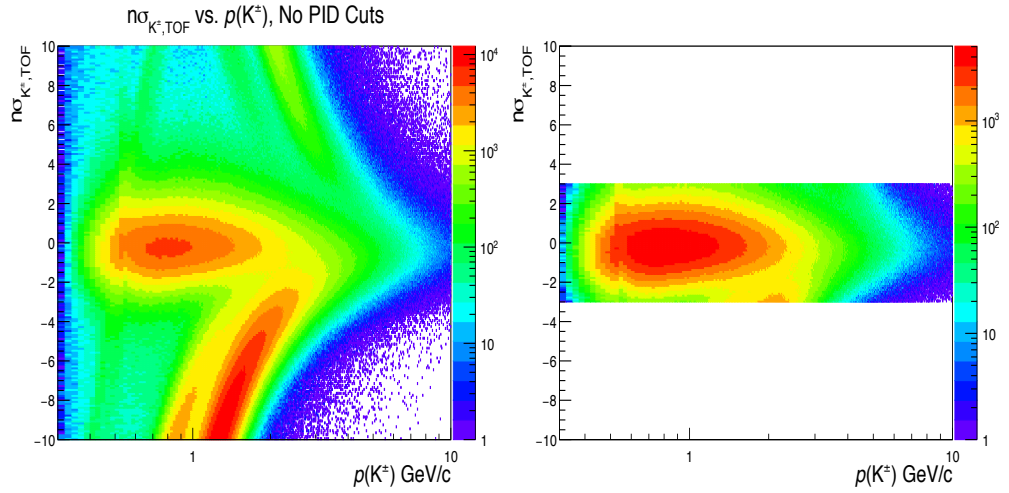


Figure 3.8: (Color online) $n\sigma_{k, TOF}$ vs. p for two different PID cut schemes. **Left Panel:** no PID cuts applied. **Right Panel:** a $|n\sigma_{k, TOF}| < 3$ cut applied [38–40].

of πK (KK) pairs is required to lie within the range $|y_{\text{pair}}| < 0.5$.

Event-mixing method is used to estimate the uncorrelated background and its shape is estimated from the invariant mass distribution of unlike-sign $K\pi$ or KK combinations from different events as shown in Fig. 3.9, 3.10, 3.11, 3.12, and 3.13. To avoid mismatch due to different acceptances and to ensure a similar event structure, tracks are mixed with events having similar vertex positions ($\Delta z < 1$ cm) and track multiplicities ($\Delta n < 5$). A reasonable statistical uncertainties are achieved by mixing events with 5 other similar events. The mixed event background is normalized in order to have the same integral as the unlike-charge distribution in the mass range $1.1 < M_{K\pi} < 1.5$ GeV/ c^2 ($1.04 < M_{KK} < 1.06$ GeV/ c^2) for $K^{*0}(\phi)$. This combinatorial background is subtracted from the unlike-charge mass distribution in each p_T bin to get the $K^{*0}(\phi)$ signal and are shown in Fig. 3.14, 3.15, 3.16, 3.17, 3.18, and 3.19. Due to an imperfect description of the combinatorial background, as well as a real correlated background, a residual background still remains after the subtraction of the combinatorial background. The real correlated background can arise from correlated $K\pi$ (KK) pairs for $K^{*0}(\phi)$, misidentified particle decays or jets.

The K^{*0} yield is extracted from the $K\pi$ invariant mass distribution in 28 p_T bins between 0 to 20 GeV/ c . After the combinatorial background subtraction, the invariant mass distribution is fitted with the combination of a Breit-Wigner function for the signal peak and a second-order polynomial for the residual background. The fit function for K^{*0} is given by

$$\frac{dN}{dM_{K^{\pm}\pi^{\mp}}} = \frac{A}{2\pi} \times \frac{\Gamma_0}{(M_{K^{\pm}\pi^{\mp}} - m_0)^2 + \frac{\Gamma_0^2}{4}} + (BM_{K^{\pm}\pi^{\mp}}^2 + CM_{K^{\pm}\pi^{\mp}} + D). \quad (3.1)$$

Here m_0 is the reconstructed mass of K^{*0} , Γ_0 is the resonance width fixed to the value in vacuum [32] and A is the yield of the K^{*0} meson. The detector mass resolution of K^{*0} is negligible compared to its width (47.4 ± 0.6 MeV/ c^2) and is therefore not included in the K^{*0} fit function. B , C and D are the fit parameters in the second order polynomial.

Similarly, the ϕ yield is extracted from the KK invariant mass distribution in 23 p_T -bins between 0.4 to 16 GeV/ c after the combinatorial background

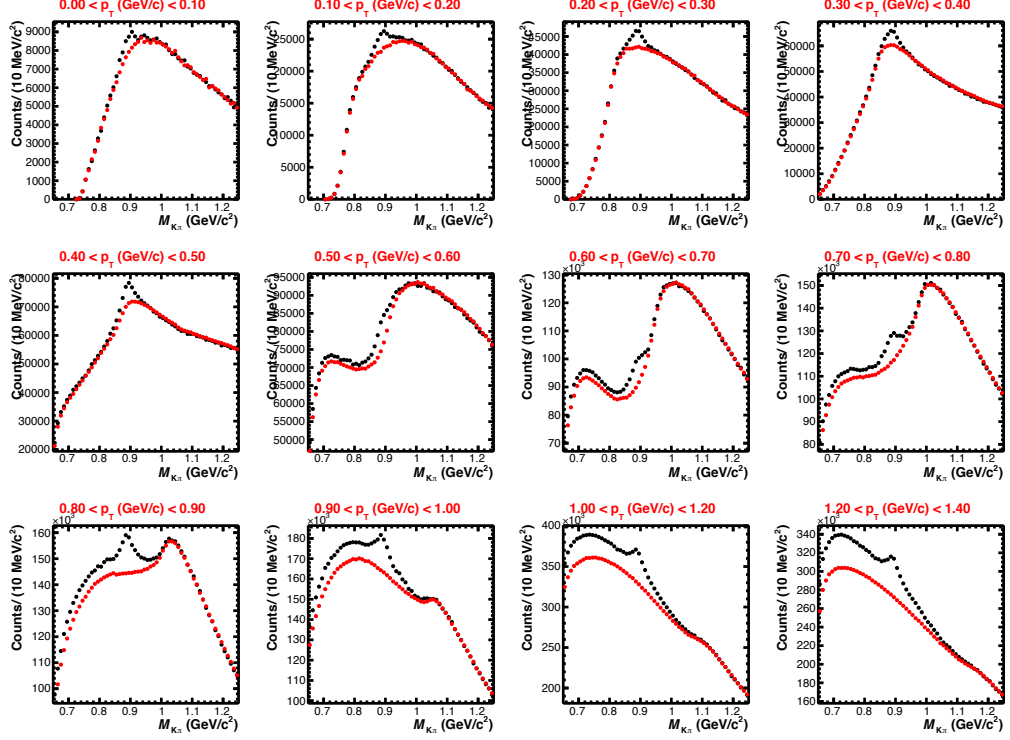


Figure 3.9: Invariant-mass distributions of K^{*0} for $0 < p_T < 1.4$ GeV/c. The black markers show the invariant-mass distribution of unlike-charge $K^\pm\pi^\mp$ pairs. Dotted red line shows the mixed event background [38–40].

subtraction. For the ϕ fit function, we consider the detector mass resolution into account due to the smaller width of the ϕ meson. This has been achieved by using a Breit-Wigner function convoluted with a Gaussian function, which is known as Voigtian function. The KK invariant mass distribution is fitted with the combination of a Voigtian function for the signal peak and a second order polynomial for the residual background. The fit function for ϕ is given by

$$\frac{dN}{dM_{KK}} = \frac{A\Gamma_0}{(2\pi^{3/2})\sigma} \int_{-\infty}^{+\infty} \exp\left(-\frac{(M_{KK} - m')^2}{2\sigma^2}\right) \frac{1}{(m' - m_0)^2 + \frac{\Gamma_0^2}{4}} dm' + (BM_{KK}^2 + CM_{KK} + D). \quad (3.2)$$

Here m_0 is the reconstructed mass of ϕ , Γ_0 is the resonance width fixed to the value in vacuum and σ is the p_T -dependent mass resolution, which ranges from 1-3 MeV/ c^2 .

To extract the yields of $K^{*0}(\phi)$, for each p_T bin we integrate the invariant mass histogram over the region $0.801 < m_{K^{*0}} < 0.990$ ($1.01 < m_\phi < 1.03$) and

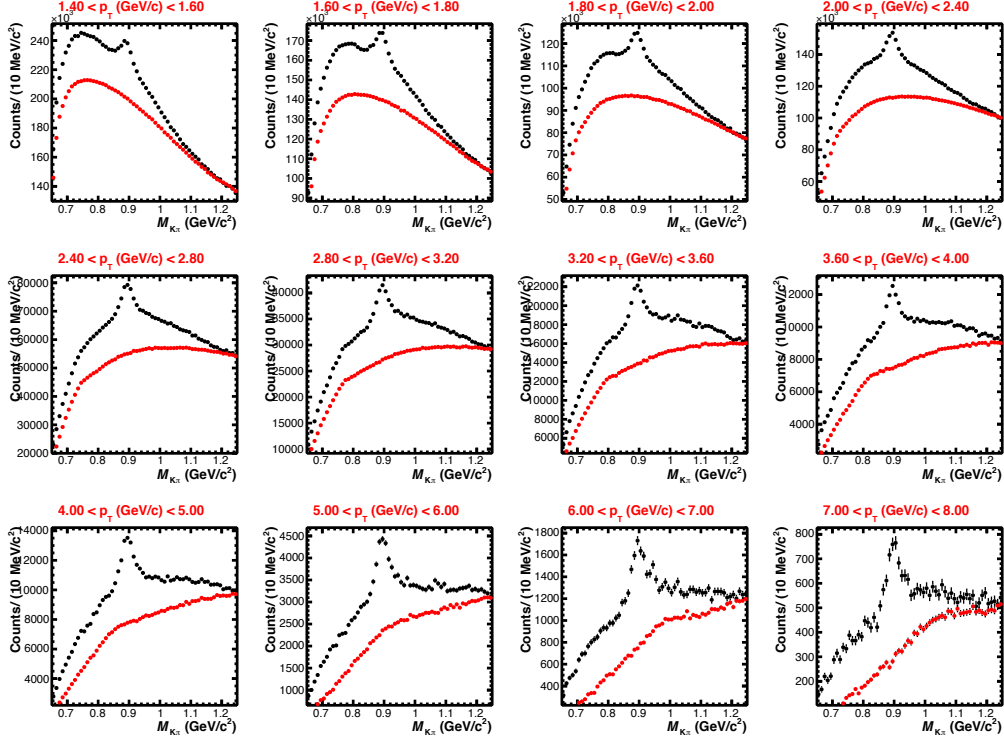


Figure 3.10: Invariant-mass distributions of K^{*0} for $1.4 < p_T < 8.0$ GeV/c. The black markers show the invariant-mass distribution of unlike-charge $K^\pm \pi^\mp$ pairs. Dotted red line shows the mixed event background [38–40].

subtract the integral of the residual background function in that same range. Raw yields are determined as:

$$N_{Raw} = N_{Bincount} - N_{Pol2} \quad (3.3)$$

where, $N_{Bincount}$ = total number of counts from the histogram in the mass range $0.801 < m_{K^{*0}} < 0.990$ ($1.01 < m_\phi < 1.03$) and N_{Pol2} = integral of the second order polynomial (Pol2) in the same range as $N_{Bincount}$.

The error in the N_{raw} is calculated by the error propagation formula, which is given as

$$\sigma_{N_{Raw}} = \sqrt{\sigma_{N_{Bincount}}^2 + \sigma_{N_{Pol2}}^2} \quad (3.4)$$

here, $\sigma_{N_{Bincount}}$ = statistical error in $N_{Bincount}$ and $\sigma_{N_{pol2}}$ = statistical error in pol2.

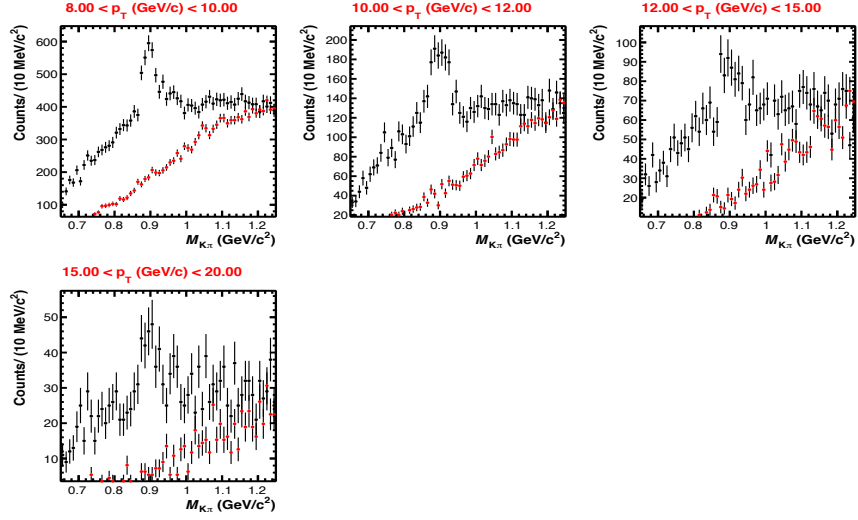


Figure 3.11: Invariant-mass distributions of K^{*0} for $8 < p_T < 20$ GeV/c. The black markers show the invariant-mass distribution of unlike-charge $K^\pm\pi^\mp$ pairs. Dotted red line shows the mixed event background [38–40].

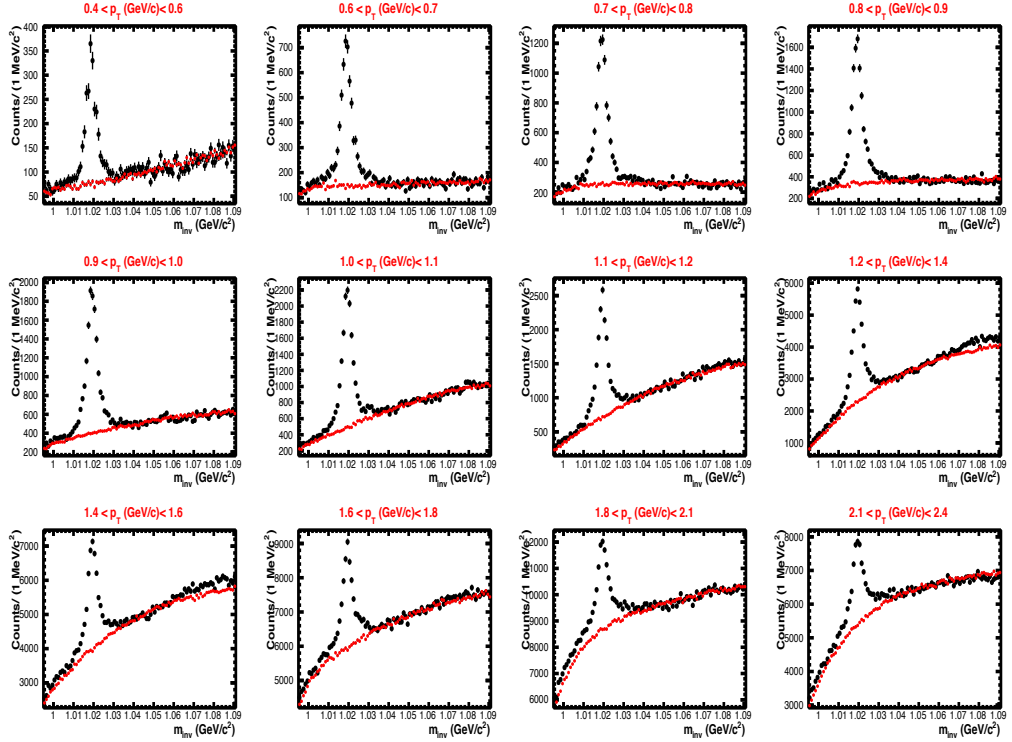


Figure 3.12: Invariant-mass distributions of ϕ for $0.4 < p_T < 16$ GeV/c for tpc2s combination. The black markers show the invariant-mass distribution of unlike-charge kaon pairs. Red marker shows the normalized event-mixing background [38–40].

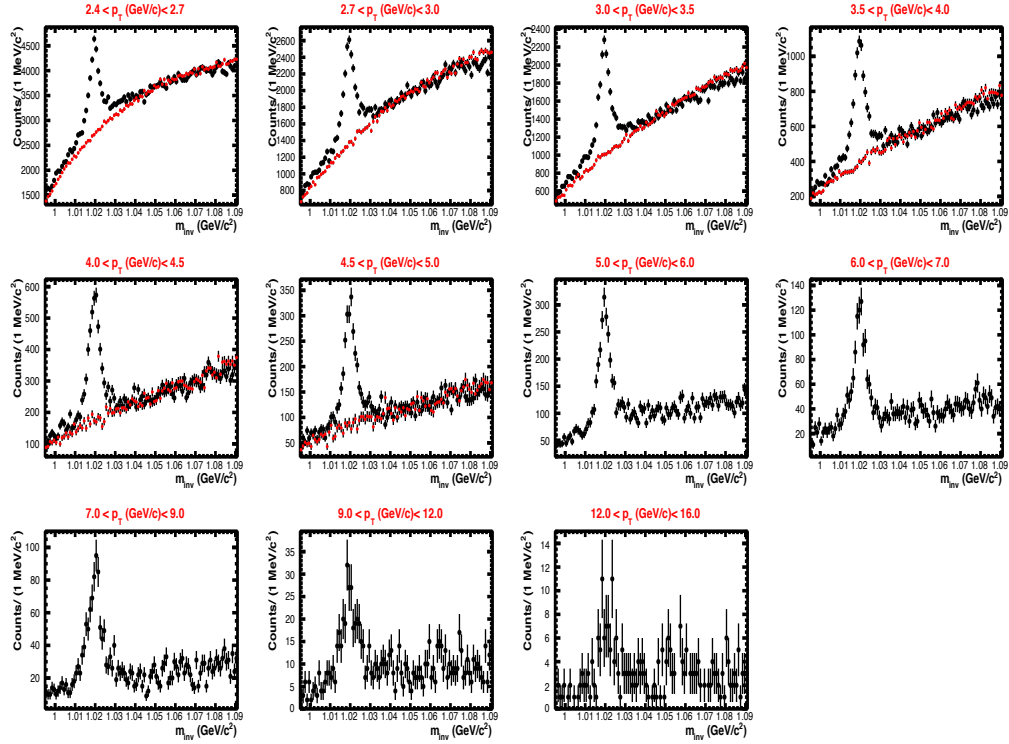


Figure 3.13: Invariant-mass distributions of ϕ for $0.4 < p_T < 16$ GeV/c for tpc2s combination. The black markers show the invariant-mass distribution of unlike-charge kaon pairs. Red marker shows the normalized event-mixing background [38–40].

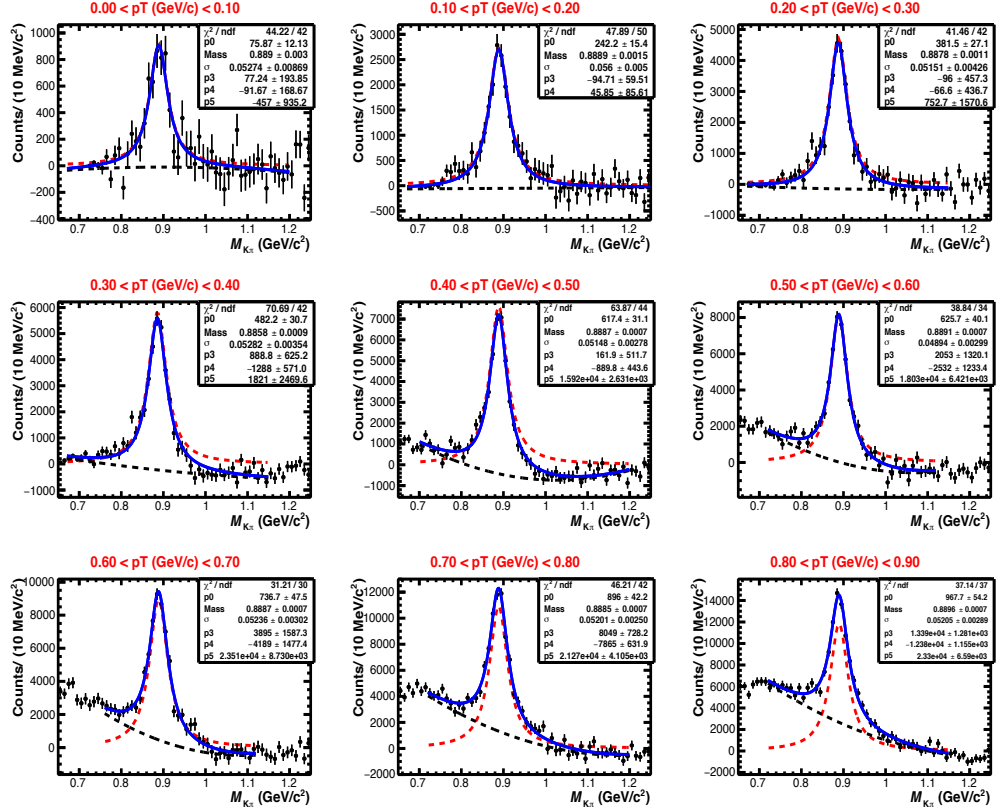


Figure 3.14: Invariant-mass distributions of K^{*0} and fits for p_T bins ranging from $0 < p_T < 0.9$ GeV/c after event-mixing background subtraction. The residual background is fitted with 2^{nd} order polynomial. Blue line shows the fitting of signal + residual background. Dotted line show the signal (red) and residual background (black) separately [38–40].

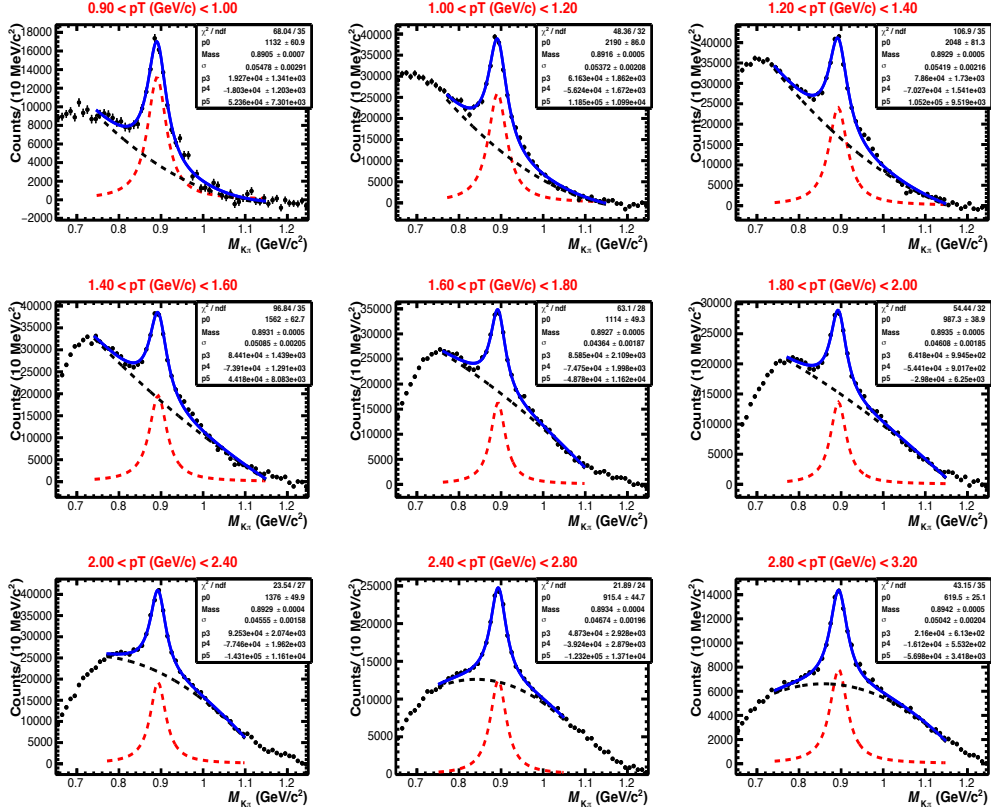


Figure 3.15: Invariant-mass distributions of K^{*0} and fits for p_T bins ranging from $0.9 < p_T < 3.2$ GeV/c after event-mixing background subtraction. The residual background is fitted with 2^{nd} order polynomial. Blue line shows the fitting of signal + residual background. Dotted line show the signal (red) and residual background (black) separately [38–40].

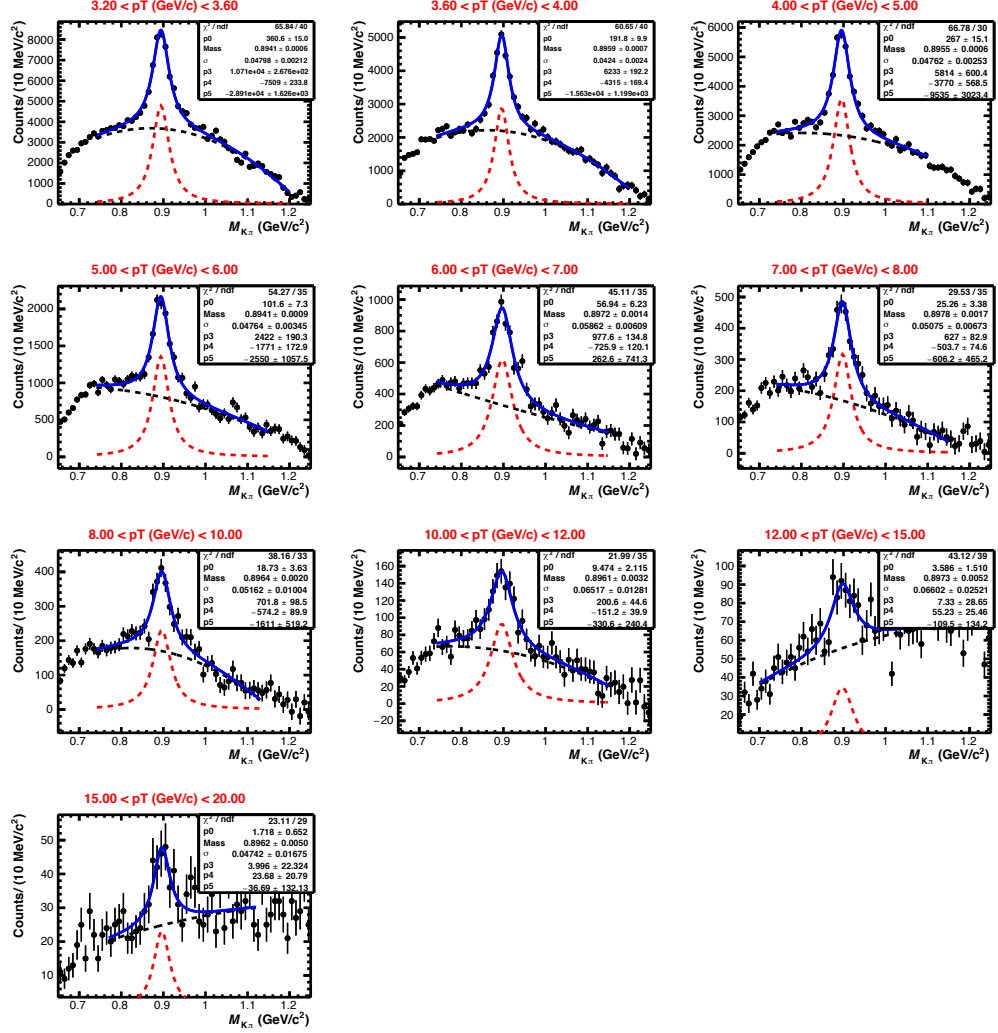


Figure 3.16: Invariant-mass distributions of K^{*0} and fits for p_T bins ranging from $3.2 < p_T < 20$ GeV/c after event-mixing background subtraction. The residual background is fitted with 2^{nd} order polynomial. Blue line shows the fitting of signal + residual background. Dotted line show the signal (red) and residual background (black) separately [38–40].

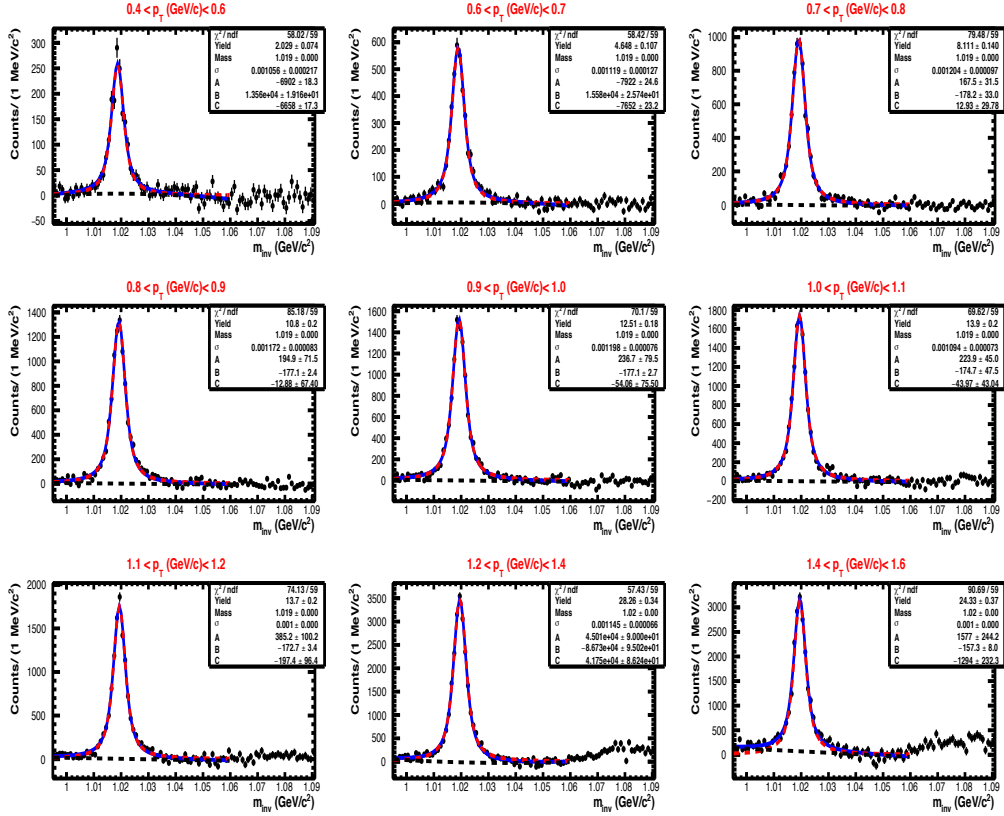


Figure 3.17: Invariant-mass distributions of ϕ and fits for p_T bins ranging from $0.4 < p_T < 1.6$ GeV/c after event-mixing background subtraction. The residual background is fitted with 2nd order polynomial. Blue line shows the fitting of signal + residual background. Dotted line show the signal (red) and residual background (black) separately [38–40].

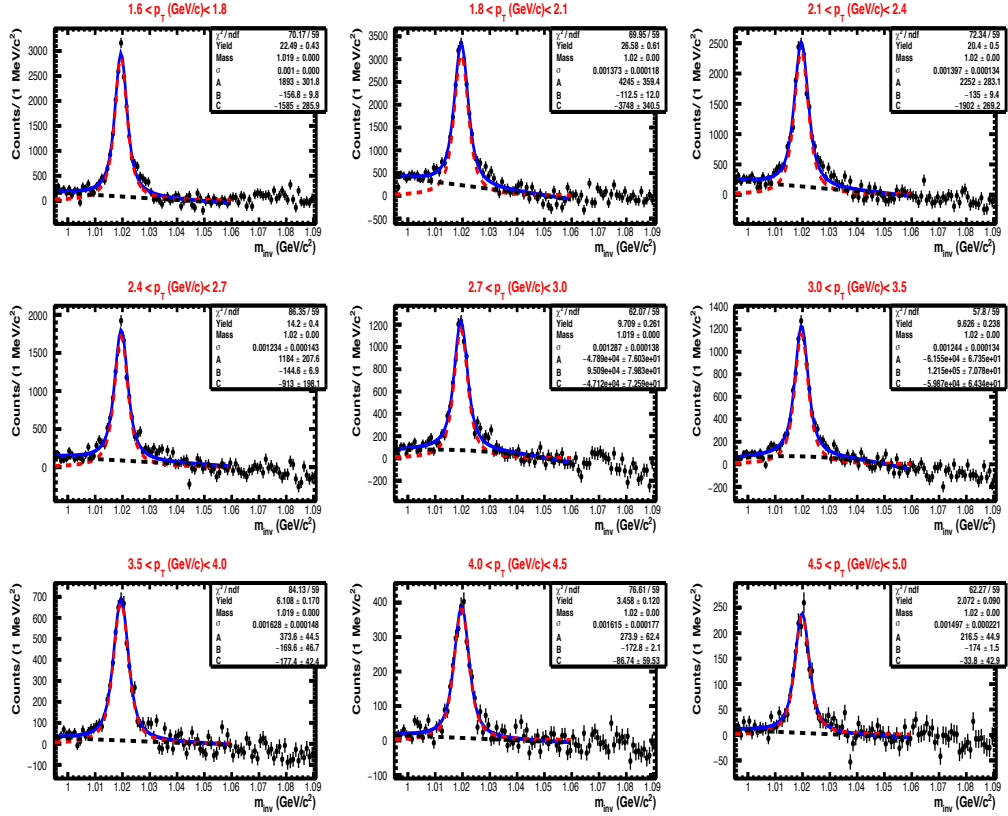


Figure 3.18: Invariant-mass distributions of ϕ and fits for p_T bins ranging from $1.6 < p_T < 5$ GeV/c after event-mixing background subtraction. The residual background is fitted with 2^{nd} order polynomial. Blue line shows the fitting of signal + residual background. Dotted line show the signal (red) and residual background (black) separately [38–40].

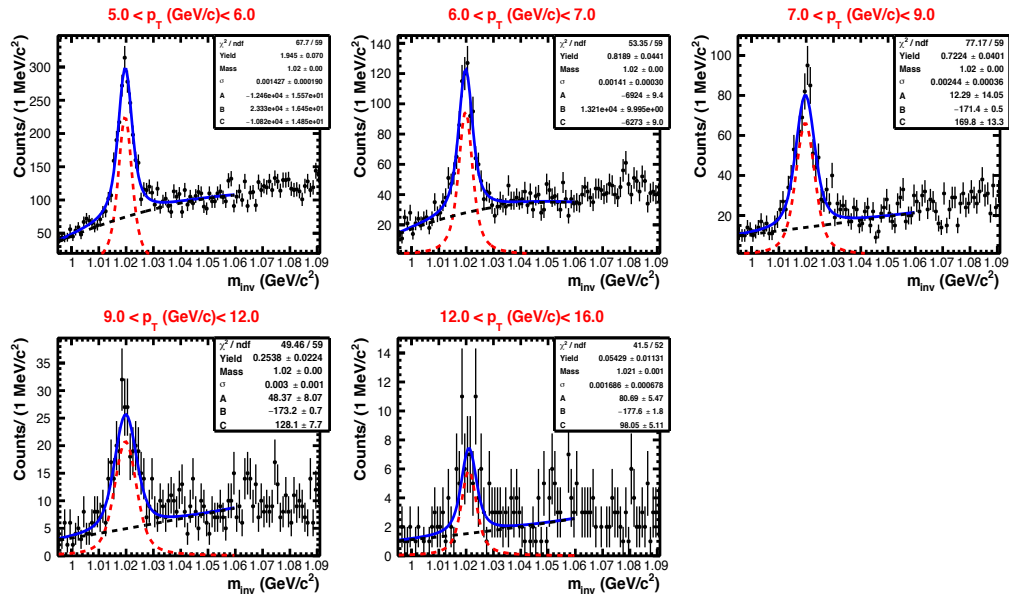


Figure 3.19: Invariant-mass distributions of ϕ and fits for p_T bins ranging from $5 < p_T < 16$ GeV/c. The residual background is fitted with 2nd order polynomial. Blue line shows the fitting of signal + residual background. Dotted line show the signal (red) and residual background (black) separately [38–40].

To take care of the tails outside the integration regions, yields are extracted from the peak fit function and added to the yields calculated from the histogram. It is assumed that the statistical uncertainties of the yields in the tail regions are fully correlated with each other and with the statistical uncertainty of the yield calculated from the histogram.

We estimate the tail contribution from the peak fit function in the mass region $(M_k + M_\pi) < m_{K^{*0}} < 0.801$ and $m_{K^{*0}} > 0.990$ for K^{*0} $\left((M_K + M_K) < m_\phi < 1.01 \text{ and } m_\phi > 1.03 \text{ for } \phi \right)$

$$N_{low} = \int_{M_K + M_\pi(M_K + M_K)}^{0.801(1.010)} f_{(signal)} dm \quad (3.5)$$

$$N_{high} = \int_{0.990(1.030)}^{\infty} f_{(signal)} dm \quad (3.6)$$

so the final raw yield is, $Y_{final} = N_{low} + N_{raw} + N_{high}$.

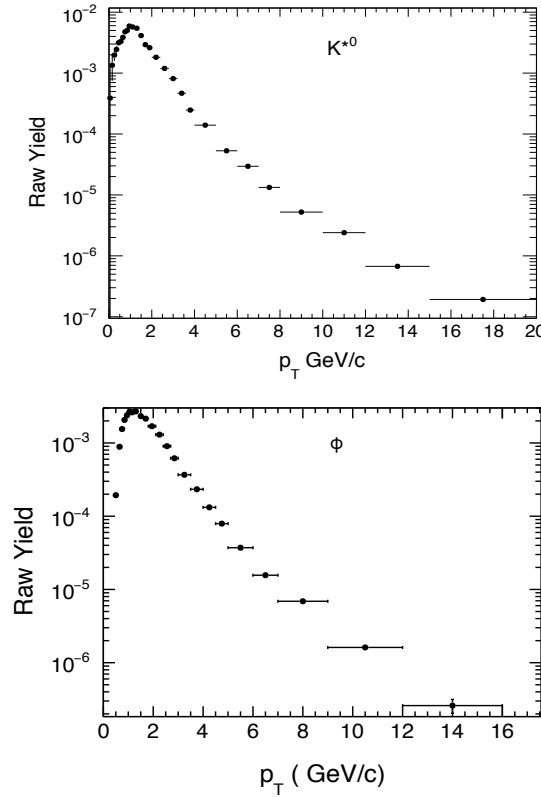


Figure 3.20: Raw yields of K^{*0} and ϕ as a function of p_T [38–40].

Figure 3.20 shows raw integrated yields as a function of p_T for K^{*0} (ϕ).

3.4 Raw Yield Normalisation and Correction

The K^{*0} and ϕ raw yields (N_{raw}) are normalised to the number of inelastic pp collisions and corrected for the hadronic branching ratio (BR), vertex correction, ALICE detector geometric acceptance (A) and efficiency (ε), and finally for signal loss. The K^{*0} and ϕ corrected yields are obtained by

$$\frac{d^2N}{dp_T dy} = \frac{N_{\text{raw}} \cdot \epsilon_{\text{SL}}}{N_{\text{evt}} \cdot \text{BR} \cdot dp_T \cdot dy \cdot \varepsilon_{\text{rec}}} \times f_{\text{norm}} \times f_{\text{vtx}}. \quad (3.7)$$

Here $\varepsilon_{\text{rec}} = A \times \varepsilon$ is the correction factor which accounts for the detector acceptance and efficiency. The ϵ_{SL} is the signal loss correction factor that accounts for the loss of $K^{*0}(\phi)$ mesons incurred by selecting events that satisfy only the ALICE minimum-bias trigger conditions, rather than all inelastic events. This is a p_T -dependent correction factor and also depends on particle species which is peaked at low p_T , indicating that events that fail the trigger selection have softer p_T spectra than the average inelastic event. This correction is the ratio of the p_T spectrum in inelastic events to the p_T spectrum in triggered events and is evaluated using Monte Carlo simulations. Let's discuss about these correction factors in details.

3.4.1 Reconstruction Efficiency $\times$ Acceptance

A simulated data set, which contains same run numbers as data is analyzed in order to extract the K^{*0} (ϕ) reconstruction efficiency $\times$ acceptance. The simulated data used for this analysis consists of about 81 million events in the the LHC15h1a1, LHC15h1b, LHC15h1c, LHC15h1d, LHC15h1h, LHC15h1i productions (AODs). Particle production and decays are simulated using PYTHIA 8 minimum bias production, while particle interactions with the ALICE detector are simulated using GEANT3. The same event selection and track quality cuts are used for the real and simulated data.

The particles produced by the event generator (without any detector effects) are referred to as the “generated” particles. These particles are the input for the GEANT3 detector simulation and the track and signal reconstruction algorithms. The tracks which are identified by the reconstruction algorithms and which pass

the track selection and PID cuts are referred to as “reconstructed” tracks.

In each transverse-momentum bin, ε_{rec} is the ratio of two quantities

Denominator: Generated $K^{*0}(\phi)$: the number of generated $K^{*0}(\phi)$ mesons with $|y_{\text{gen}}| < 0.5$ that decayed to pair of $K^\pm\pi^\mp$ (K^+K^-).

Numerator: Reconstructed $K^{*0}(\phi)$: the number of $K^{*0}(\phi)$ mesons with $|y_{\text{rec}}| < 0.5$ for which $K^\pm\pi^\mp$ (K^+K^-) decay daughters are reconstructed.

The ε_{rec} as a function of p_T is shown in Figure 3.21. It has a strong dependence on transverse momentum at low p_T . The uncertainty in ε_{rec} is calculated using the Bayesian approach [33]. The standard deviation in an efficiency $\varepsilon = k/n$, where the numerator k is a subset of the denominator n , is

$$\sigma_\varepsilon = \sqrt{\frac{k+1}{n+2} \left(\frac{k+2}{n+3} - \frac{k+1}{n+2} \right)}. \quad (3.8)$$

The fractional statistical uncertainty in ε_{rec} is added in quadrature with the statistical uncertainty in the uncorrected $K^{*0}(\phi)$ yield to give the total statistical uncertainty of the corrected $K^{*0}(\phi)$ yield.

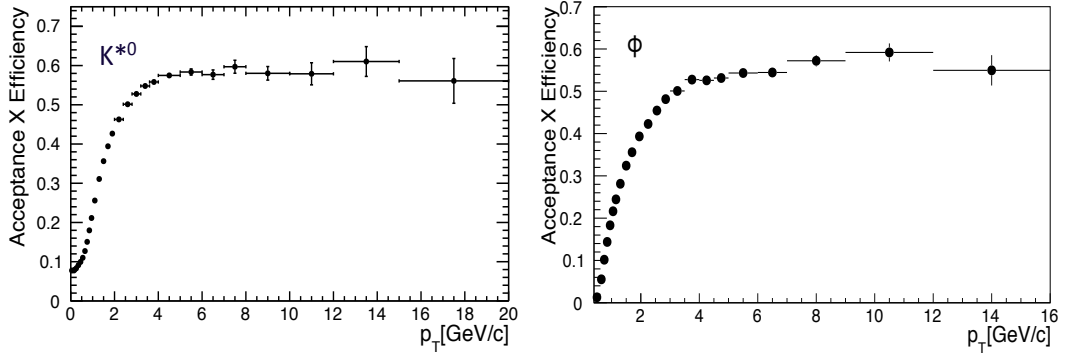


Figure 3.21: $K^{*0}(\phi)$ meson reconstruction efficiency (ε_{rec}) as a function of p_T for tpc2stofv3s [TPC 2σ with TOF veto 3σ cut] (tpc2s [TPC 2σ cut]) [38–40].

The efficiency for the combined period is calculated by taking the weighted average of efficiencies of individual data periods.

$$Efficiency_{All \text{ data periods}}(\varepsilon_{\text{rec}}) = \sum_i w_i \varepsilon_i(p_T) \quad (3.9)$$

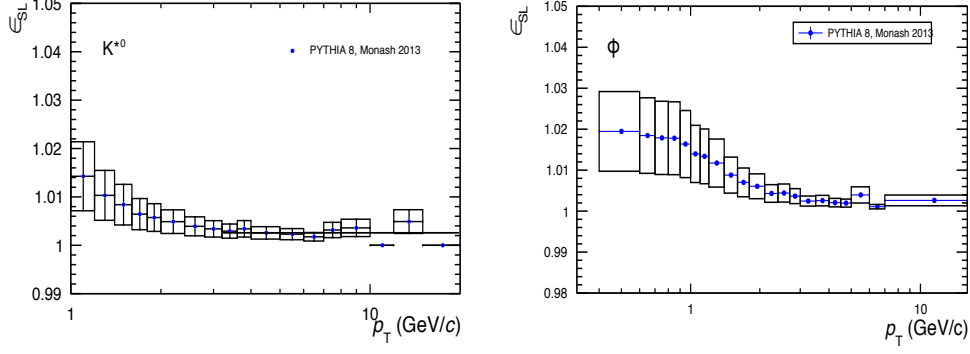


Figure 3.22: Signal loss correction factor as a function of p_T . The boxes represent the systematic uncertainties [38–40].

$$\sigma_{\epsilon_{rec}} = \sqrt{\sum_i w_i^2 \sigma_{\epsilon_i}^2(p_T)} \quad (3.10)$$

Here $w_i = \frac{\text{Accepted events in data period}_i}{\text{Total accepted events}}$ and $i = \text{LHC12a, LHC12b, LHC12c, LHC12d, LHC12h, LHC12i}$.

3.4.2 Signal-Loss Correction

The signal loss correction factor is calculated to account for the loss of K^{*0} (ϕ) mesons incurred by selecting events that satisfy only kINT7 trigger, rather than all inelastic events. This is a p_T -dependent correction factor which is peaked at low p_T , indicating that events that fail the kINT7 selection have softer K^{*0} (ϕ) p_T spectra than the average inelastic event. This correction is the ratio of the p_T spectrum in inelastic events to the p_T spectrum in kINT7 events.

The correction factor is calculated using PYTHIA 8 simulation production. The signal-loss as a function of p_T is shown in the Fig. 3.22. The uncertainty in the figure is $[0.5 \times |\epsilon_{SL} - 1|]$.

3.4.3 Inelastic Normalization Factor

In order to normalize to the number of inelastic pp collisions, the inelastic normalization factor (f_{norm}) is calculated. This is needed to convert a particle yield normalized to the number of triggered events to a yield normalized to the number of inelastic events. The value of inelastic normalization factor is 0.7718 ± 0.0228 [42].

3.4.4 f_{vtx} Correction Factor

The factor f_{vtx} is also needed in the normalization and it is a ratio with

Denominator: the number of triggered events (after application of the pileup cuts)

Numerator: the subset of the events in the denominator for which a good vertex was found (*i.e.*, it passes the vertex quality cuts, but without the cut on the z -position of the vertex)

The vertex correction factor for K^{*0} analysis is 0.972 and for ϕ this is calculated separately for each period and then the mean value is taken as final value which is 0.978.

3.5 Estimation of Systematic Uncertainties

The systematic uncertainties are due to different sources such as signal extraction, background subtraction, track selection, global tracking uncertainty, knowledge of the material budget and the hadronic interaction cross-section as summarised in Table 3.1. The calculations of the systematic uncertainties for K^{*0} and ϕ are similar to those for our previous analyses [41].

Systematic Uncertainty for K^{*0}

For each transverse-momentum bin there are many measurements of the yield and its statistical uncertainty. In general, the yield is calculated for every possible permutation of the analysis parameters (e.g. fit region variations, background subtraction variation, normalization range variations, PID cuts, track cuts.....etc). The general strategy for evaluating systematic uncertainties is:

- Choose one set of analysis parameters as the “default” set.
- Observe the deviations in the yield from the default measurement when one of those parameters is changed.
- The RMS value is taken to be the systematic uncertainty for a given source.

Table 3.1: Systematic uncertainties in the measurement of K^{*0} and ϕ yields in pp collisions at $\sqrt{s} = 8$ TeV. The global tracking uncertainty is p_T -independent, while the other single-valued systematic uncertainties are averaged over p_T . The values given in ranges are minimum and maximum uncertainties depending on p_T .

Source	K^{*0} (%)	ϕ (%)
Signal extraction	7.1	1.5
Background subtraction	5.0	1.2
Track selection	4.0	2.0
Material budget	0 - 3.4	0 - 3.1
Hadronic Interaction	0 - 2.8	0 - 5.4
Global tracking efficiency	6.0	6.0
Branching ratio	–	1.0
Total	11.3 - 12.1	6.7 - 9.1

- The systematic uncertainties for different sources are added in quadrature to obtain the total systematic uncertainty from all sources.

List of sources considered for the systematic study are as follows

1. Source of Uncertainties

- (a) Signal Extraction
- (b) Background Subtraction method
- (c) Track and Event Selection

2. Ungrouped uncertainties

- (a) Material Budget
- (b) Hadronic interaction cross-section
- (c) Signal Loss uncertainty
- (d) Inelastic Normalization

Let us denote the yield in the default case as y_{def} with statistical error σ_{def} and yield in the systematic case as y_{sys} with statistical error σ_{sys} .

The default set of parameters chosen for K^{*0} measurements are given below:

- **Combinatorial background:** Mixed event
- **Normalization Region:** Area normalization out side the K^{*0} meson peak in the region $1.1 < m_{inv} < 1.5 \text{ GeV}/c^2$.
- **Fit Region:** Depending up on the background shape the fit region is chosen for each p_T bin for default case.
- **Residual Background Fit:** Quadratic polynomial (*pol2*).
- **Peak Fit:** : Breit-Wigner, where the width parameter is a free parameter.
- **Yield Extraction Method:** Bin counting.
- **PID Cut:** Both Pions and Kaons are selected by a cut of $|N\sigma_{\text{TPC}}| < 2.0$ with a TOF veto of $|N\sigma_{\text{TOF}}| < 3.0$. It should be noted that the value quoted for the TPC PID cut above only apply for $p(K^\pm) > 0.4 \text{ GeV}/c$ and $p(\pi^\pm) > 0.4 \text{ GeV}/c$. Wider TPC PID cuts are used at low momentum due to problems with TPC PID. Specifically, for $p(K^\pm) < 0.3 \text{ GeV}/c$ and $p(\pi^\pm) < 0.3 \text{ GeV}/c$, a cut of $|N\sigma_{\text{TPC}}| < 6.0$ is applied for both Kaons and Pions, while a cut of $|N\sigma_{\text{TPC}}| < 4.0$ for $0.3 < p < 0.4 \text{ GeV}/c$ is applied for both Kaons and Pions.

1. Signal extraction

Fit range variation: The following three fitting ranges are chosen for the systematic study.

- (a) Fit region 0 (sh: Fit 0): $0.76\text{-}1.12 \text{ GeV}/c$
- (b) Fit region 1 (sh: Fit 1): $0.73\text{-}1.12 \text{ GeV}/c$
- (c) Fit region 2 (sh: Fit 2): $0.77\text{-}1.14 \text{ GeV}/c$

The ratio to default is shown in the Fig. 3.23

Normalization range variation: The following three normalisation ranges are chosen for the systematic study. The ratio to default is shown in the Fig. 3.24.

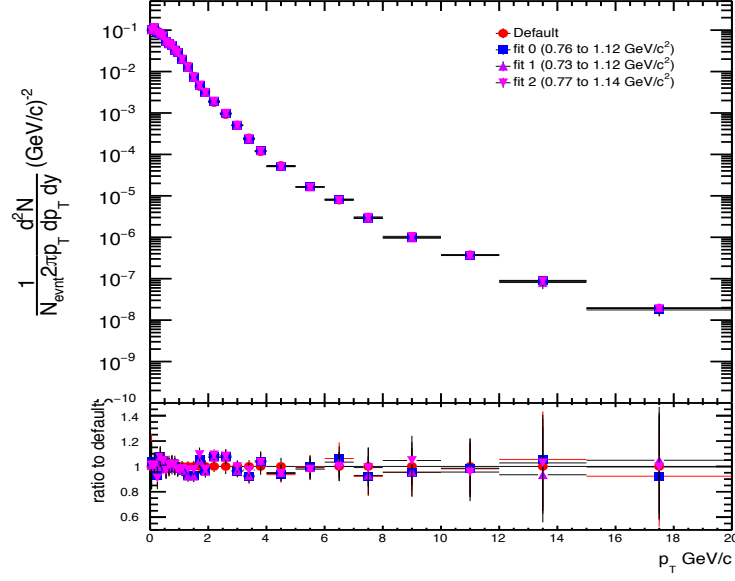


Figure 3.23: Corrected K^{*0} spectrum and ratios to default due to fit range variations [38–40].

- (a) Normalisation region 0 (sh: Mix 0): 1.1–1.3 GeV/c
- (b) Normalisation region 1 (sh: Mix 1): 1.1–1.4 GeV/c
- (c) Normalisation region 2 (sh: Mix 2): 1.15–1.4 GeV/c

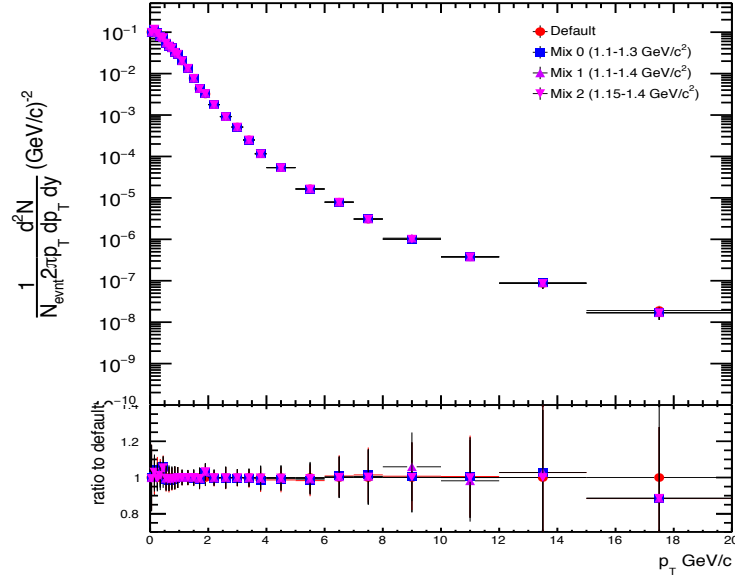


Figure 3.24: Corrected K^{*0} spectrum and ratios to default due to normalisation range variations [38–40].

Residual background function variation: For systematic study, we

have taken the polynomial of order three (pol3) as residual background variation.

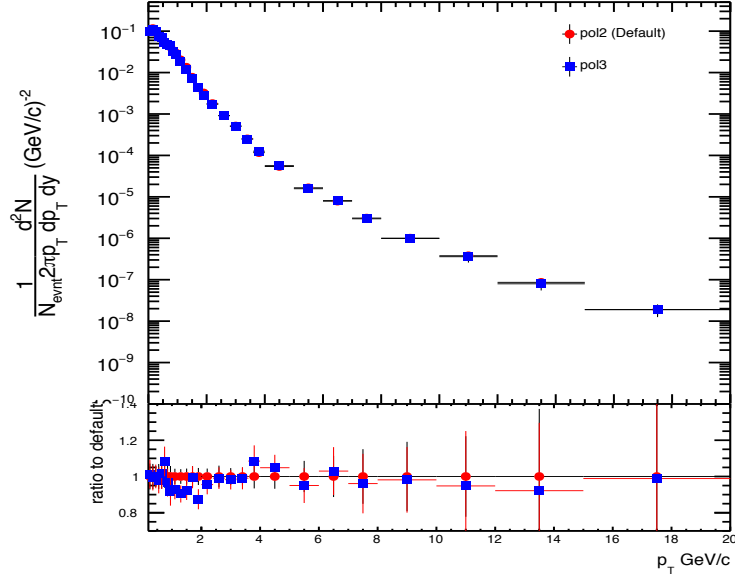


Figure 3.25: Corrected K^{*0} spectrum and ratios to default due to residual background variations [38–40].

The ratio to default is shown in the Fig. 3.25.

Function Integration: For the raw yield extraction, we have used the bin counting as the default method. The function integration for raw yield extraction is used for the systematic study.

The ratio to default is shown in the Fig. 3.26.

Fixed Width: For the raw yield extraction, we have used the free width during fitting as the default method and fixed width to the PDG value ($0.0474 \text{ GeV}/c^2$) for systematic study. The ratio to default is shown in the Fig. 3.27.

2. **Combinatorial Backgrounds:** Combinatorial background estimated from like sign method is used as a systematic study. The ratio between raw yield due to like sign and mixed event background estimation is shown in the Fig. 3.28.

3. **Track and Event Selection:** Systematic error is calculated by varying one track cut at a time. Track cut variations are listed below. Ratios w.r.t.

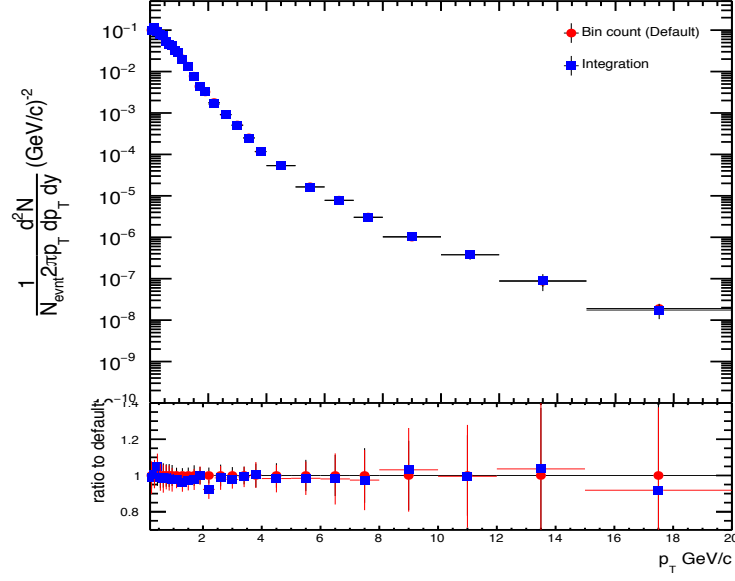


Figure 3.26: Corrected K^{*0} spectrum and ratios to default due to function integration [38–40].

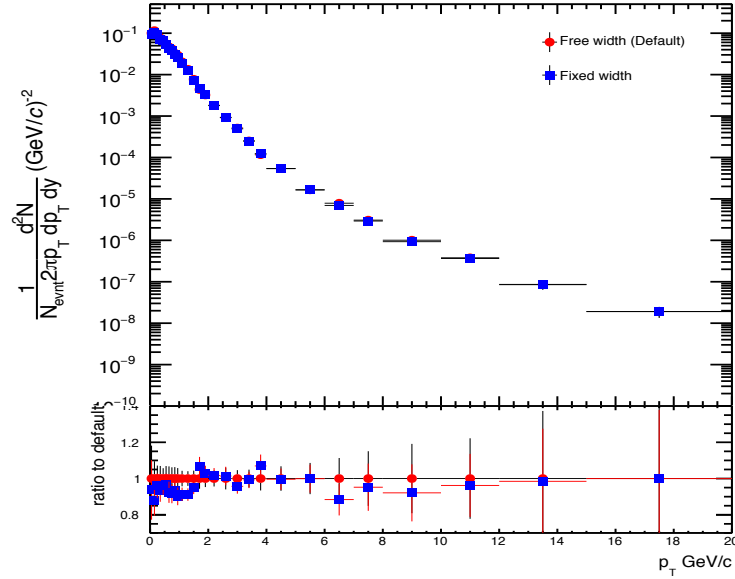


Figure 3.27: Corrected K^{*0} spectrum and ratios to default due to fixed width raw yield extraction [38–40].

default value are plotted in the lower panel.

DCA_{xy} Cut: $DCA_{xy} < 7\sigma$ is taken as default and $DCA_{xy} < 4\sigma$ is taken for systematic study. Ratio between p_T spectra due to this two track selection is given in Fig. 3.29

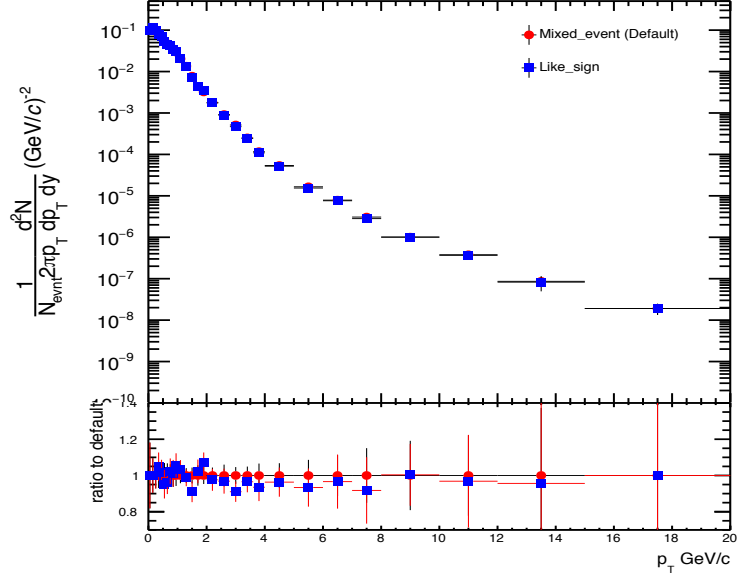


Figure 3.28: Corrected K^{*0} spectrum and ratios to default due to like sign background subtraction [38–40].

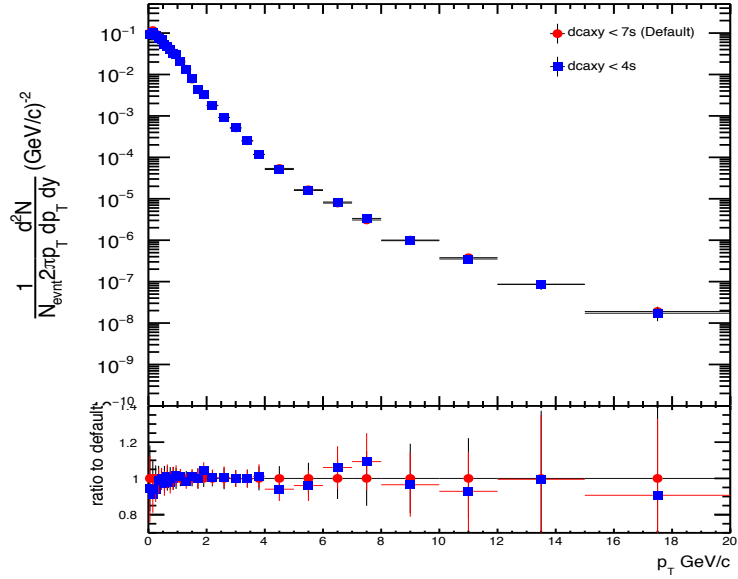


Figure 3.29: Corrected K^{*0} spectrum and ratios to default due to DCAxy cut [38–40].

DCA_z Cut: $DCA_z < 2$ cm is taken as default and $DCA_z < 1$ cm is taken for systematic study. Ratio between p_T spectra due to this two track selection is given in Fig. 3.30

Crossed-rows Cut: Number of crossed rows in TPC > 70 is taken as default and number of crossed rows in TPC > 80 and TPC > 100 are

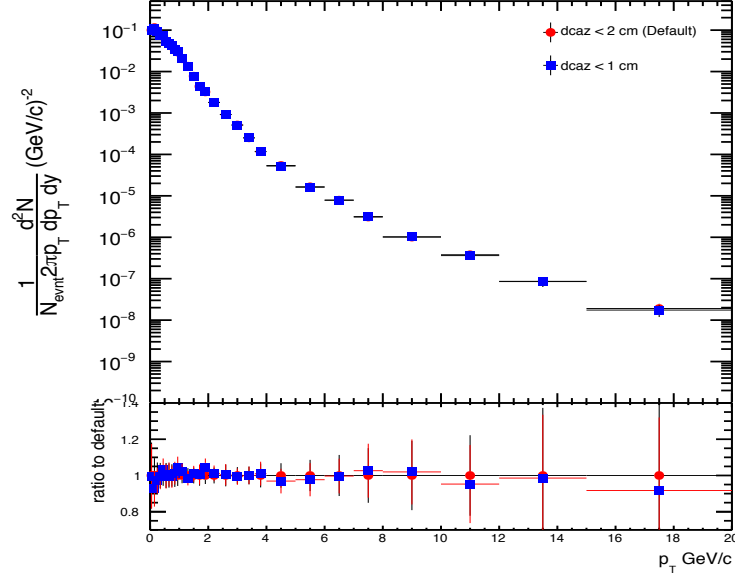


Figure 3.30: Corrected K^{*0} spectrum and ratios to default due to DCAz cut [38–40].

taken for systematic studies. Ratio between p_T spectra due to this two track selection is given in Fig. 3.31

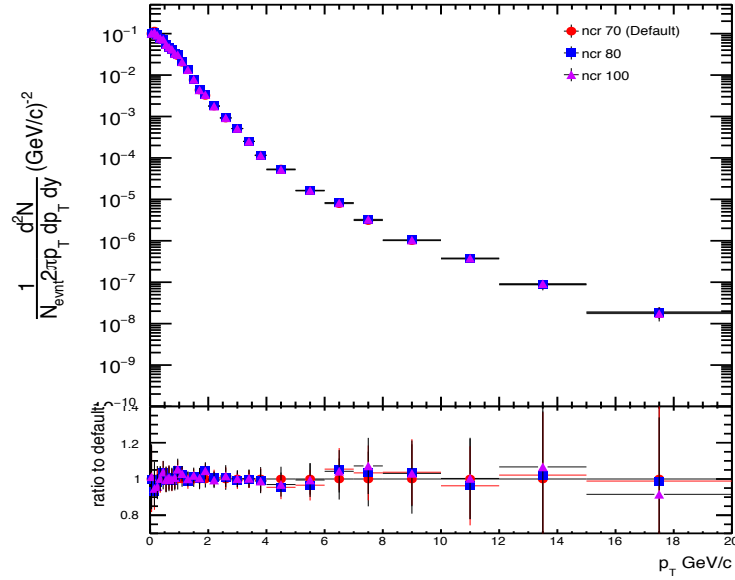


Figure 3.31: Corrected K^{*0} spectrum and ratios to default due to ncr (Number of crossed rows) cut [38–40].

Crossed rows over findable cluster Cut: Crossed rows over findable cluster < 0.8 as default and 0.9 is taken as a systematic study. Ratio

between p_T spectra due to this two track selection is given in Fig. 3.32.

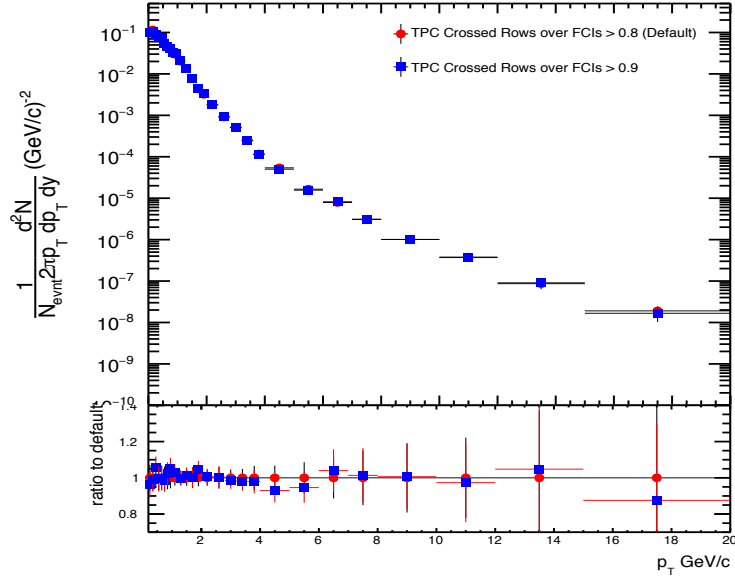


Figure 3.32: Corrected K^{*0} spectrum and ratios to default due to crossed rows over findable cluster cut [38–40].

Vertex Cut: $|v_z| < 10$ cm is taken as default. $|v_z| < 12$ cm and $|v_z| < 8$ cm is taken for systematic study. Ratio between p_T spectra due to this two event selection is given in Fig. 3.33

4. PID Selection Cut: Default : TPC 2σ TOFVETO 3σ

For systematic study : TPC 2σ TOFVETO 4σ , TPC 2.5σ TOFVETO 3σ and TPC 1.5σ TOFVETO 3σ corrected spectra for this PID cuts and their ratios are shown in Fig. 3.34.

Systematic Uncertainty for ϕ

In general, the analysis is performed many times, using every possible permutation of analysis parameters (PID cuts, combinatorial background, etc.). Broadly, the uncertainties are calculated in 2 different groups.

1. Grouped Uncertainties

(a) Signal Extraction

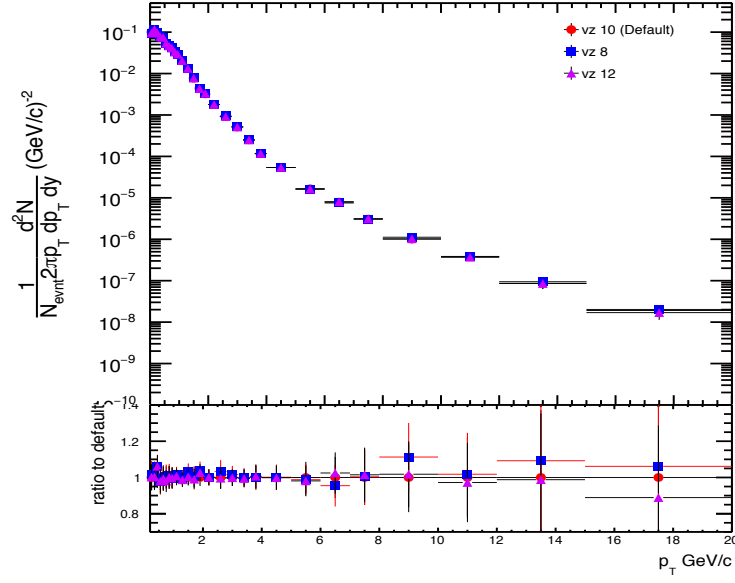


Figure 3.33: Corrected K^{*0} spectrum and ratios to default due to vertex cut [38–40].

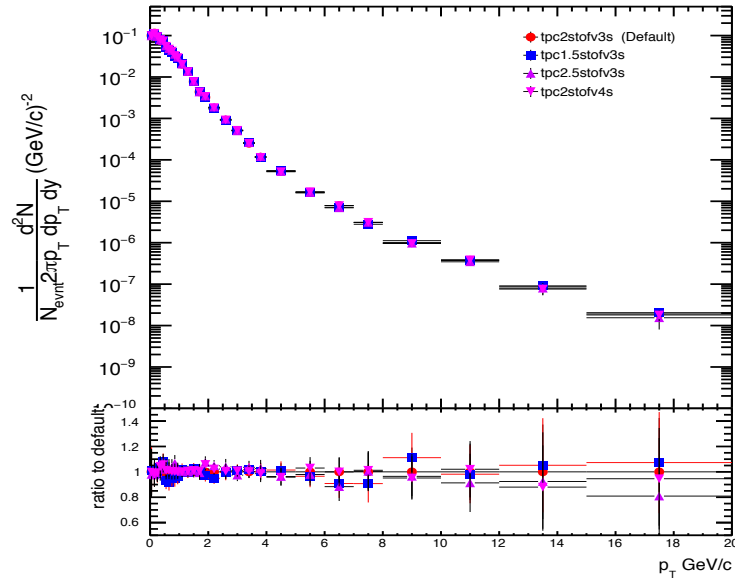


Figure 3.34: Corrected K^{*0} spectrum and ratios to default due to pid selection cut [38–40].

(b) Background Subtraction method

(c) Track Selection Cuts

2. Ungrouped uncertainties

(a) Material Budget

- (b) Hadronic interaction cross-section
- (c) Inelastic Normalization
- (d) Branching Ratio

This means that there are multiple measurements of the yield for each transverse momentum bin as follows.

1. Signal Extraction:

Event Mixing Normalization regions: For the event mixing normalization region, we have 3 choices: `mix1` ($0.995 < m < 1.005$), `mix2` ($1.04 < m < 1.06$), and `mix3` ($1.05 < m < 1.08$).

Fit Region: For fitting the invariant-mass distribution function, we have 5 choicers: `fr0` ($0.995 < m < 1.06$), `fr1` ($0.995 < m < 1.07$), `fr2` ($0.99 < m < 1.09$), `fr3` ($0.995 < m < 1.09$) and `fr4` ($0.99 < m < 1.06$).

Residual Background Fit: For residual background fit, we have 3 choices: linear (`pol1`), quadratic (`pol2`), or cubic (`pol3`)

Peak Fit: 1 choice: `voigtG`

For each background subtraction method and track selection cut, the above signal extraction variations are calculated. For these set of variations the mean value is calculated. The standard deviation from the set of measurements is the signal extraction systematic uncertainty. Then the mean statistical uncertainty is obtained. After this step, the mean value is the yield and mean statistical uncertainty is the statistical error for the particular set of background subtraction method and track selection cut. This procedure is repeated for each track selection cut and background subtraction method. Figure 3.35 shows an example of different variations for signal extraction grouping procedure.

2. Background Subtraction:

For background subtraction, we have 2 choices: like-charge (`like_sign`), or mixed event (`event_mix`). For a particular track selection cut, the mean values from the previous grouping are used. Now, for each track

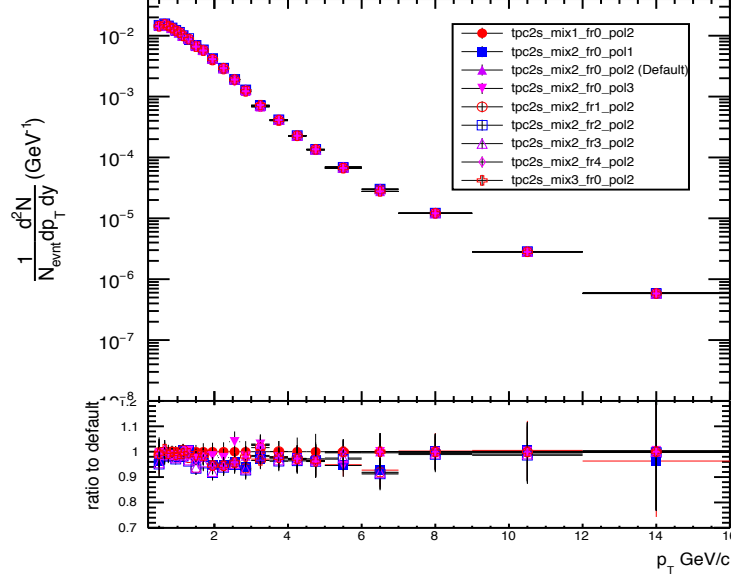


Figure 3.35: **Upper plot:** Yield from **tpc2s** after varying different signal extraction techniques, **Lower plot:** Ratio to default combination [38–40].

selection cut, background subtraction techniques are varied. So, there will be 2 measurements for each set. The standard deviation from each set is taken as the background subtraction systematic uncertainty. The mean value is the yield and mean statistical uncertainty is the statistical error for the particular set of track selection cut, and PID strategy. Figure 3.36 show the example of different variations for background subtraction.

3. Track Selection Cuts:

PID Cuts: We have 5 different choices for PID variations: $2\sigma_{K,TPC}$ with $3\sigma_{K,TOF}$ veto (**tpc2stofv3s**), $1.5\sigma_{K,TPC}$ with $3\sigma_{K,TOF}$ veto (**tpc15stofv3s**), $2.5\sigma_{K,TPC}$ with $3\sigma_{K,TOF}$ veto (**tpc25stofv3s**), $2\sigma_{K,TPC}$ (**tpc2s**), or $2\sigma_{K,TPC}$ with $4\sigma_{K,TOF}$ veto (**tpc2stofv4s**). For this analysis we consider TPC 2σ cut for PID.

Event/Track Selection Cuts: We have 10 different choices for track cut variations: event $|v_z| < 8$ cm (**vz8**), event $|v_z| < 12$ cm (**vz12**), $N_{cr,TPC} \geq 100$ (**ncr100**), $N_{cr,TPC} \geq 80$ (**ncr80**), $R_{TPC} \geq 0.9$ (**rtpc9**), $DCA_{xy} < 4\sigma$ (**dcaxy4s**), $DCA_z < 1$ cm (**dcaz1**), $\chi^2_{track} < 3.5$ (short-hand: **x2track35**)

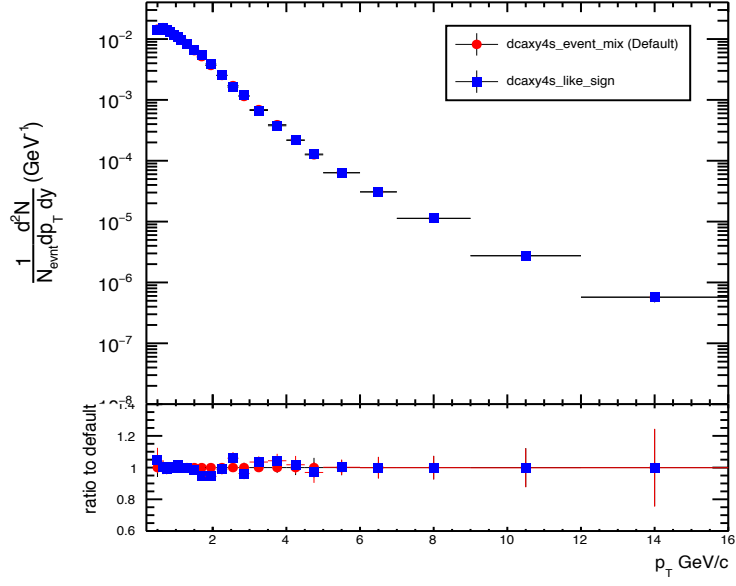


Figure 3.36: **Upper plot:** Yield of dcaxy4s after varying different background subtraction techniques, **Lower plot:** Ratio to default combination [38–40].

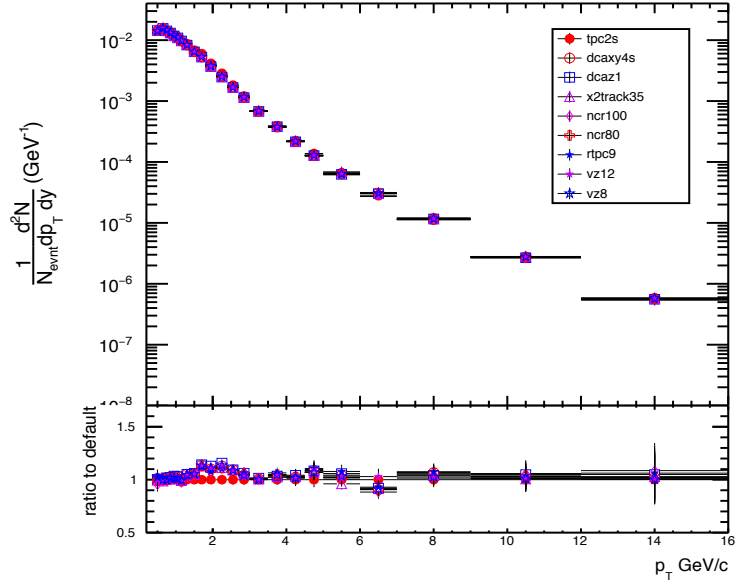


Figure 3.37: **Upper plot:** Yield from different track cuts, **Lower plot:** Ratio of all track cuts to the default [38–40].

There will be 5 TPC/TOF cuts and 9 quality cuts variation in this step. The standard deviation from these set of measurements is taken as track selection systematic uncertainty. The mean value is the yield and mean statistical uncertainty is the final statistical error. Figure 3.37 shows the yield from different PID cuts, event-selection cuts and track quality cuts.

Now lets discuss about the ungrouped systematic uncertainties considered for K^{*0} and ϕ .

1. **Material Budget:** The systematic uncertainty due to uncertainties in the ALICE material budget is taken from the analysis of the K^{*0} (ϕ) mesons in p-Pb collisions [43–45]. It is p_T -dependent and reaches a maximum value of a few percent.
2. **Hadronic Interaction Cross-section:** The systematic uncertainty due to uncertainties in the ALICE hadronic interaction cross-section is taken from the analysis of the K^{*0} (ϕ) mesons in p-Pb collisions [43–45]. It is p_T -dependent and reaches a maximum value of a few percent.
3. **Inelastic Normalization:** The systematic uncertainty due to the uncertainty in the factor f_{norm} (used to normalize to the number of inelastic pp collisions) has a value of 2.95%.
4. **Branching Ratio:** The systematic uncertainty due to the uncertainty in the branching ratio is 1% for ϕ meson and independent of p_T .
5. **Tracking Uncertainty:** The values of tracking uncertainty for 8 TeV is taken as 3% each for daughter particles [46]. So the tracking uncertainty for K^{*0} (ϕ) mesons is 6%.

The total systematic uncertainty is the sum in quadrature of the uncertainties from the sources listed above.

3.5.1 Consistency Checks

These analyses are repeated many times by changing the analysis parameters to estimate the systematic uncertainty. However, we must make sure that the sys-

tematics arising due to these variations are not due to statistical fluctuations. As described in [47], the condition for consistency is *not* agreement within the statistical uncertainties: since the default and alternate measurements use (mostly) overlapping data sets, agreement within the statistical uncertainties is not sufficient. Instead, agreement within the difference in quadrature of the statistical uncertainties is required. If two measurements are found to be consistent within the quadrature difference of their statistical uncertainties, then it is not necessary to account for the differences between them in the systematic uncertainties.

Let us denote the yield in the default case as y_{def} with statistical error σ_{def} and yield in the alternate case as y_{alt} with statistical error, σ_{alt} . The difference between the yields is denoted as $\Delta = y_{def} - y_{alt}$ and the quadrature difference of their statistical error is $\sigma = \sqrt{\sigma_{def}^2 + \sigma_{alt}^2}$. Then we define a factor $n = \Delta/\sigma$, where n is calculated for each p_T bin and a distribution of n or Δ/σ is plotted in Fig. 3.38 and 3.39. In general, if two measurements are consistent, it is expected that the distribution of n would have a mean near 0, a standard deviation near 1, and that 68% (95%) of the entries would lie within $|\Delta/\sigma| < 1$ ($|\Delta/\sigma| < 2$). For this analysis we consider a source as a systematic source if $|\Delta/\sigma|$ distribution of that source having standard deviation > 1 . Therefore, it appears that the measurements are not consistent with the default measurement; it is necessary to account for the differences between these measurements in the systematic uncertainties. This method is called Barlow check.

Figure 3.40 shows the fractional uncertainties from all the sources as a function of p_T .

3.6 Results and Discussion

3.6.1 Transverse Momentum Spectra and Differential p_T Ratio

In this analysis we measure the production of K^{*0} and ϕ in pp collisions at $\sqrt{s} = 8$ TeV in the range $0 < p_T < 20$ GeV/ c for K^{*0} and $0.4 < p_T < 16$ GeV/ c for ϕ . Fig. 3.41 shows the corrected transverse momentum spectra of K^{*0} and ϕ at midrapidity $|y| < 0.5$ after taking into account all sources of uncertainties.

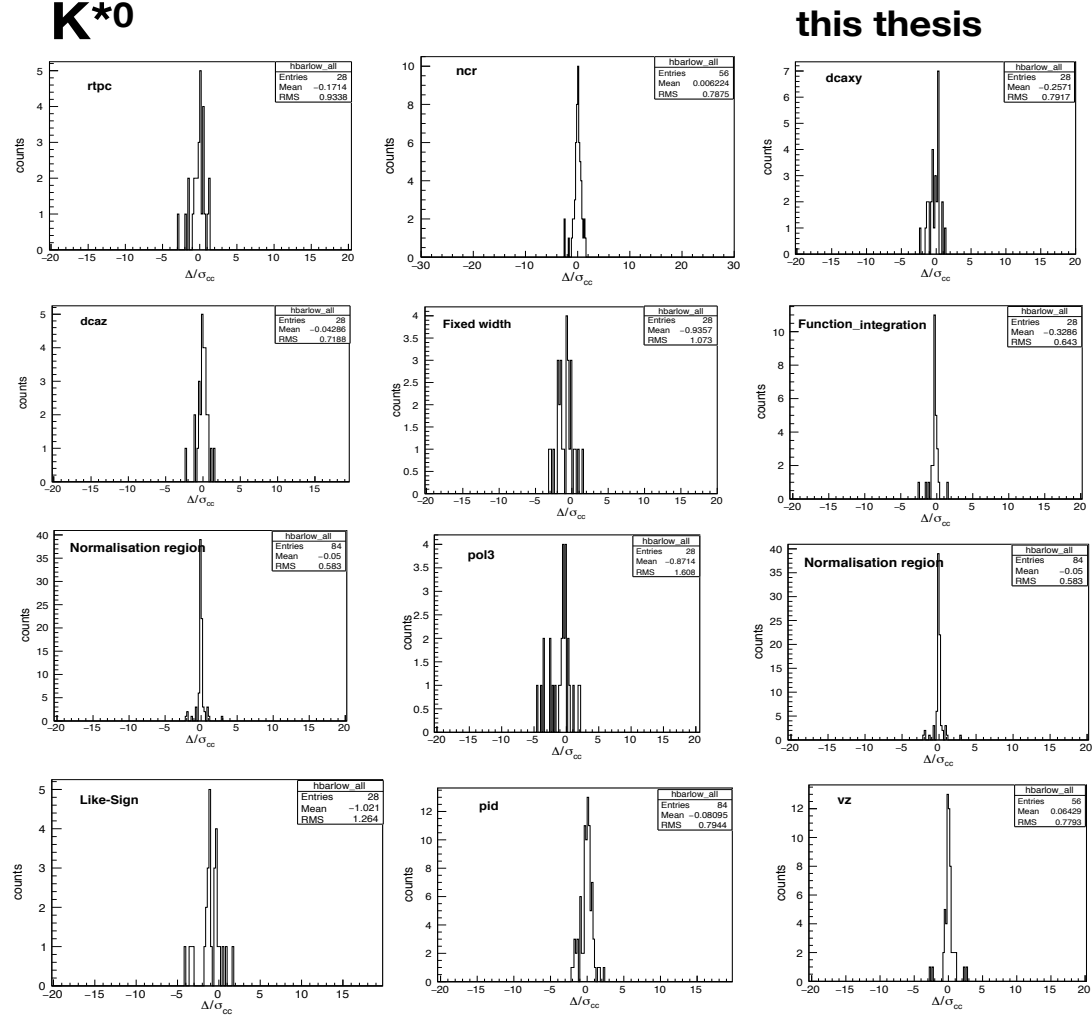


Figure 3.38: Consistency checks for variations in signal extraction techniques for K^{*0} meson [38–40].

We have studied the energy evolution of the transverse momentum spectra for K^{*0} and ϕ by taking the ratio of differential yields for different LHC energies to the yields at $\sqrt{s} = 2.76$ TeV and this is shown in Fig 3.42. It manifests from Fig. 3.42 that the differential yield ratio is independent of p_T within the systematic uncertainties up to 1 GeV/c for different collision energies. This suggests that the particle production mechanism in soft scattering regions is independent of collision energy. An increase in slope of differential yield ratios is observed for $p_T > 1$ GeV/c, indicating a larger number of hard scatterings as the collision energy increases.

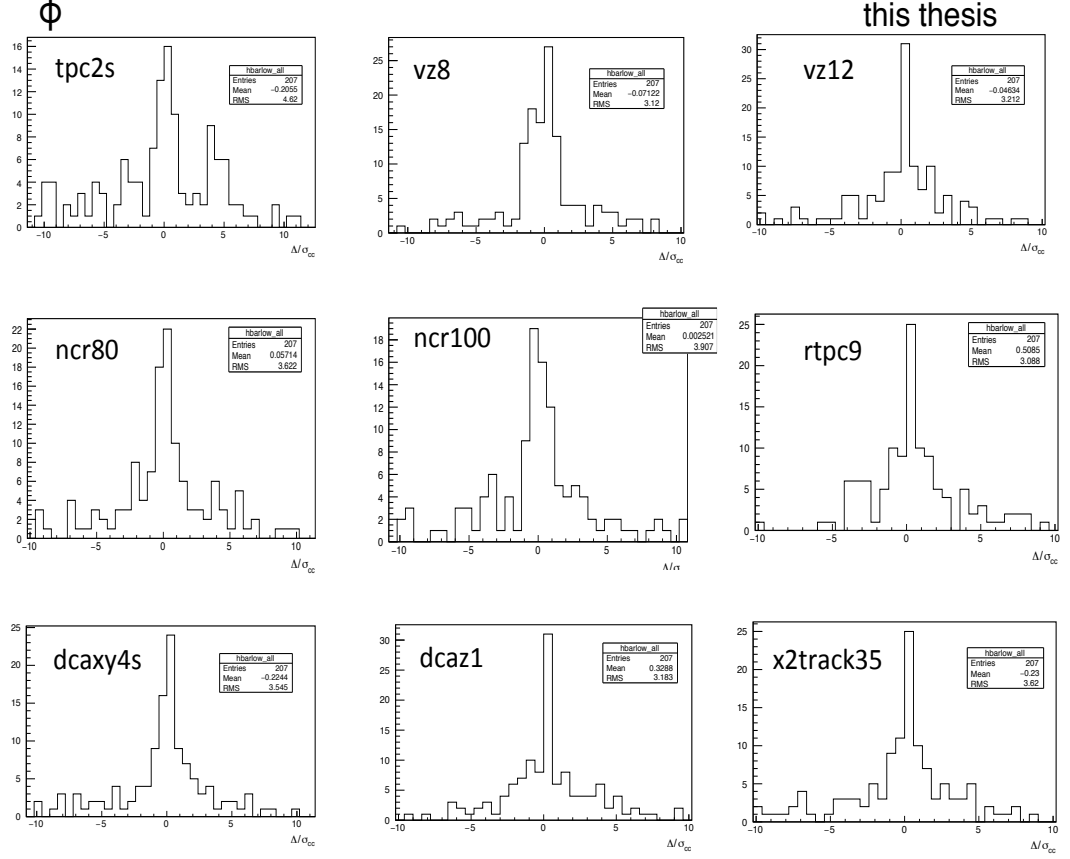


Figure 3.39: Consistency checks for variations in signal extraction techniques for ϕ meson [38–40].

3.6.2 Integrated Yields and $\langle p_T \rangle$

Table 3.2 shows the K^{*0} and ϕ integrated yields (dN/dy) and mean transverse momenta ($\langle p_T \rangle$) in minimum bias pp collisions at $\sqrt{s} = 8$ TeV. As the ϕ spectrum starts from 0.4 GeV/c, for the calculation of integrated yields and $\langle p_T \rangle$, the spectrum is extrapolated to 0 GeV/c using a Lévy-Tsallis fit [48, 49].

Table 3.2: K^{*0} and ϕ integrated yields and $\langle p_T \rangle$ in minimum bias pp collisions at $\sqrt{s} = 8$ TeV.

Particles	measured p_T (GeV/c)	dN/dy	$\langle p_T \rangle$ (GeV/c)
K^{*0}	0.0 - 20.0	$0.101 \pm 0.001 \pm 0.014$	$1.037 \pm 0.006 \pm 0.029$
ϕ	0.4 - 16.0	$0.0335 \pm 0.0003 \pm 0.0030$	$1.146 \pm 0.005 \pm 0.040$

This section describes the procedure used to fit the ϕ -meson p_T spectrum

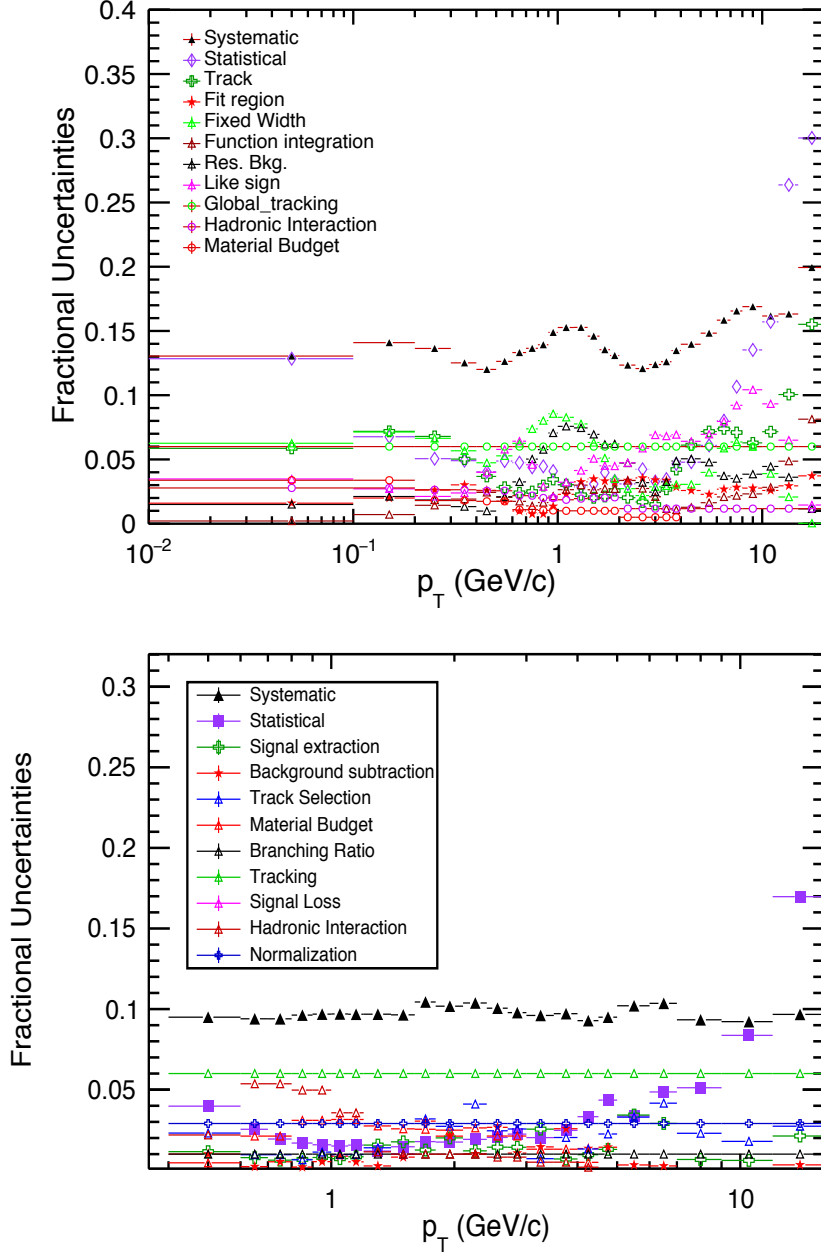


Figure 3.40: (Color online) Fractional uncertainties from all the sources for K^{*0} (top panel) and ϕ (bottom panel) [38–40].

and the quantities derived from fits of the spectrum.

Procedure for Extrapolation

The normalized, corrected ϕ meson spectrum ($d^2N/dp_T dy$) is fit with several different functions, which always include a multiplicative pre-factor of p_T .

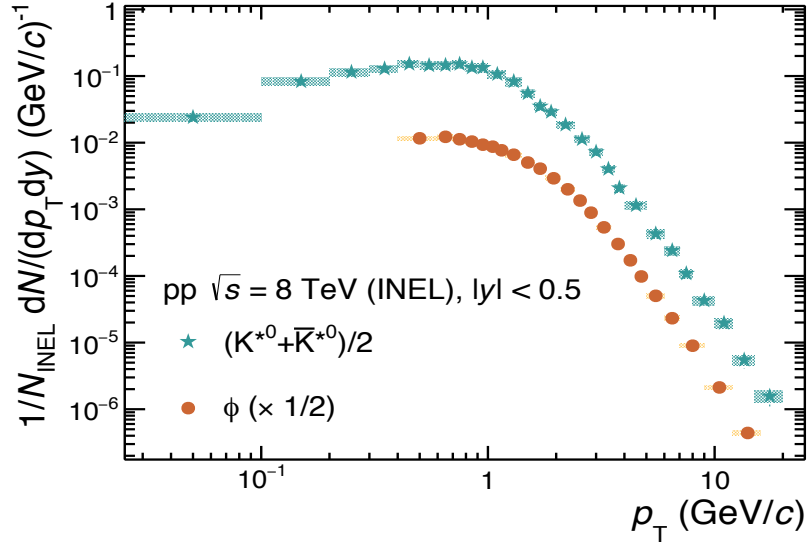


Figure 3.41: (Color online) p_T spectra of K^{*0} and ϕ in inelastic pp collisions at 8 TeV. The bars (usually not visible) are statistical uncertainties and the boxes are systematic uncertainties [38–40].

1. **An exponential function in m_T :**

$$f(p_T) = p_T \times A \times (e^{-m_T/T}) \times \frac{e^{M/T}}{T^2 + TM}. \quad (3.11)$$

Here, M is the mass of the ϕ meson, $m_T = \sqrt{M^2 + p_T^2}$ and the function is normalized so that the scale parameter A gives its integral. For this fit function, the histogram is only fit over the range $0.4 < p_T < 2$ GeV/c, giving a better fit in the low- p_T region. This is included as a systematic uncertainty.

2. **An exponential function in p_T :**

$$f(p_T) = p_T \times A \times (e^{-p_T/T}) \times \frac{e^{M/T}}{T^2 + TM}. \quad (3.12)$$

Here, M is the mass of the ϕ meson and the function is normalized so that the scale parameter A gives its integral. For this fit function, the histogram is only fit over the range $0.4 < p_T < 2$ GeV/c. This is included as a systematic uncertainty.

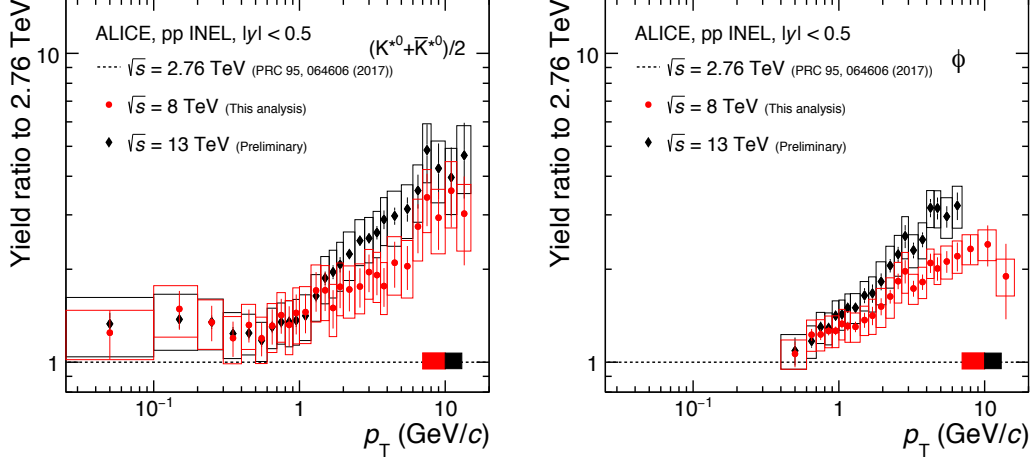


Figure 3.42: (Color online) Ratios of transverse-momentum spectra of K^{*0} and ϕ in INEL events at various center-of-mass energies to the spectra in pp collisions at 2.76 TeV. The statistical and systematic uncertainties are shown as vertical error bars and boxes, respectively. The normalisation uncertainties are indicated by boxes around unity [38–40].

3. A Boltzmann distribution:

$$f(p_T) = p_T \times A \times m_T \times (e^{-m_T/T}) \times \frac{e^{M/T}}{2T^3 + 2T^2M + TM^2}. \quad (3.13)$$

As with the closely related m_T exponential, this function is normalized so that the scale parameter A gives its integral and the fit is only performed over the range $0.4 < p_T < 2$ GeV/c. This is included as a systematic uncertainty.

4. **Lévy-Tsallis function:** The histogram is fit over the range $0.4 < p_T < 2.5$ GeV/c using a Lévy-Tsallis function [48, 49]. This is used as the **default function**:

$$f(p_T) = p_T \times \frac{A(n-1)(n-2)}{nT[nT + M(n-2)]} \left[1 + \frac{m_T - M}{nT} \right]^{-n}. \quad (3.14)$$

The Lévy-Tsallis function fit is also done for different p_T ranges such as: $0.4 < p_T < 3$ GeV/c and $0.4 < p_T < 3.5$ GeV/c and they are included for the calculation of systematic uncertainty in integrated yield and mean transverse momentum.

5. **Bose-Einstein and Fermi-Dirac distribution:** The histogram is fit over the range $0.4 < p_T < 2$ GeV/c using a Bose-Einstein distribution :

$$f(p_T) = p_T \times A \times \frac{1}{e^{m_T/T} - 1} \times (e^{M/T} - 1), \quad (3.15)$$

These functions are normalized so that when the pre-factor of p_T is removed, each has a value of A at $p_T = 0$. These fittings are included as systematic uncertainty.

6. **Power-law function:** The histogram is fit for $2 < p_T < 16$ GeV/c with a power-law function :

$$f(p_T) = p_T \times A \times (1 + p_T/T)^{-n}. \quad (3.16)$$

Table 3.3: χ^2/ndf and fitting parameters for fits

Functions	Range (GeV/c)	χ^2/ndf	A	n	T (GeV)
Levy-Tsallis	0.4-2.5	0.0624	0.0348 ± 0.002	4.585 ± 1.174	0.255 ± 0.036
Levy-Tsallis	0.4-3	0.150	0.033 ± 0.001	5.769 ± 1.029	0.281 ± 0.027
Levy-Tsallis	0.4-16	0.2591	0.0328 ± 0.001	7.271 ± 0.203	0.306 ± 0.008
m_T -exponential	0.4-2	0.9911	0.030 ± 0.001	—	0.386 ± 0.0138
p_T -exponential	0.4-2	0.0887	0.014 ± 0.001	—	0.518 ± 0.0165
Boltzmann	0.4-2	0.9911	0.045 ± 0.001	—	0.387 ± 0.014
Bose-Einstein	0.4-2	0.858	0.142 ± 0.009	—	0.395 ± 0.0151
Power-law	2-16	0.143	1.693 ± 1.291	7.257 ± 0.429	1.437 ± 0.359

Figure 3.43 shows the fitting of listed functions. Each fit is performed twice. First, each spectrum is fit with the statistical and systematic uncertainty for each point added in quadrature; collectively, this set of fits will be denoted as $f(p_T)$. Second, each spectrum is fit using only its statistical uncertainties; collectively, this set of fits will be denoted $f'(p_T)$. The set f is the primary set of fits and will be used to extract dN/dy and $\langle p_T \rangle$. The set f' is used in the calculation of the statistical uncertainties of those quantities.

dN/dy for ϕ meson is calculated by integrating the histogram (which spans the range $0.4 < p_T < 16$ GeV/c) and using the spectrum fit $f(p_T)$ described in

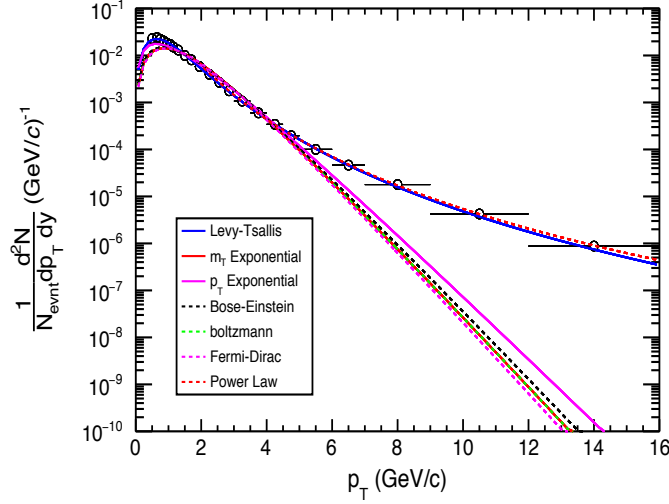


Figure 3.43: Corrected ϕ spectrum ($d^2N/dp_T dy$) with fits over the range $0 < p_T < 16$ GeV/c [38–40].

Section 3.6.2 to extrapolate the yield to the low and high- p_T regions outside the measured range. The central value is calculated using Lévy-Tsallis function which spans the range $0.4 < p_T < 2.5$ GeV/c and the function is extrapolated to 0 GeV/c and 25 GeV/c. The extrapolated region accounts for 14.8% of dN/dy at 0-0.4 GeV/c and 0.02% at 16-25 GeV/c. Low- and high- p_T extrapolation histograms with narrow bins (0.05 GeV/c wide) are generated using the spectrum fit f' . The statistical uncertainties of the bin contents of these histograms are determined using ROOT's `TF1::IntegralError()` function.

A large number (1000) of alternate histograms is randomly generated by shifting the points in the measured and extrapolation histograms. A random factor k is generated from a Gaussian with mean 0 and $\sigma = 1$. Each of the measured bins is shifted independently (*i.e.*, a different k factor is generated for each bin). All of the bins in the low- p_T extrapolation histogram are shifted coherently using the same k factor, and similarly for the high- p_T extrapolation histogram. For each iteration, the three alternate (random) histograms are integrated and the total dN/dy is stored. The mean (μ_{alt}) and standard deviation (σ_{alt}) of the set of 1000 alternate values are found. The statistical uncertainty in the default value

of dN/dy is

$$\sigma(dN/dy) = (dN/dy) \times \frac{\sigma_{\text{alt}}}{\mu_{\text{alt}}}. \quad (3.17)$$

The extrapolated part amounts to $\simeq 15\%$ of the yield. The systematic uncertainties include the contributions from the choice of the spectrum fit function for extrapolation, which amounts to $\sim 1.5\%$. In case of K^{*0} , no extrapolation is needed as the spectrum starts from 0 GeV/ c .

The values of dN/dy for $\pi^\pm$ and $K^\pm$ in pp collisions at $\sqrt{s} = 8$ TeV are interpolated using the data points available at different LHC collision energies [50, 51] namely 0.9, 7 and 13 TeV. The values at $\sqrt{s} = 8$ TeV are interpolated by fitting the data points with the following polynomial function

$$A(\sqrt{s})^n + B. \quad (3.18)$$

Here A , n and B are the fitting parameters. In order to calculate uncertainties on the interpolated value, the central values of the data points are shifted within their uncertainties and fitted with the same function. The $\pi^+ + \pi^-$ and $K^+ + K^-$ interpolated yields in minimum bias pp collisions at $\sqrt{s} = 8$ TeV are (4.54 ± 0.22) and (0.586 ± 0.027) , respectively.

3.6.3 Particle ratios

Figure. 3.44 shows the ratio of the dN/dy of K^{*0} (ϕ) to that of π in left (right) panel, as a function of the collision energy in pp collisions. Here, π has no strangeness content, K^{*0} has one unit of strangeness, and ϕ is strangeness neutral but contains two strange (anti)quarks. It is observed that the K^{*0}/π and ϕ/π ratios are independent of the collision energy within the systematic uncertainty, which indicates that the chemistry of the system is independent of energy from RHIC to LHC energies. This also suggests that the strangeness production mechanisms do not depend on energy in minimum bias pp collisions at LHC energies. Fig. 3.44 and Ref. [41] show that this flat behaviour is observed from RHIC to LHC energies and the new result at $\sqrt{s} = 8$ TeV is in agreement with previous findings. We want to stress that this flat behaviour is not trivial: since particle yields do in fact increase with the collision energy, the flat ratios are indicative

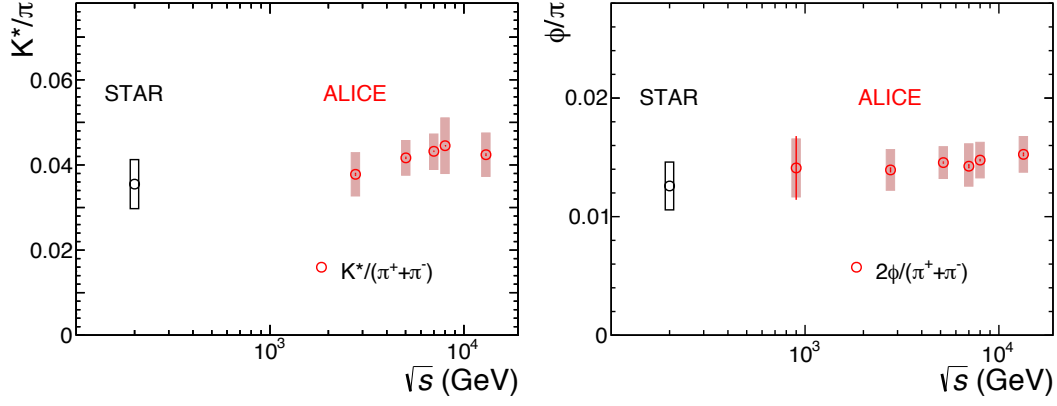


Figure 3.44: (Color online) Particle ratios of K^{*0}/π (left) and ϕ/π (right) are presented for pp collisions as a function of the collision energy. Bars (when present) represent statistical uncertainties. Boxes represent the total systematic uncertainties or the total uncertainties for cases when separate statistical uncertainties were not reported [38–40].

of the fact that the percentage increase of dN/dy for π , K^{*0} and ϕ as a function of the collision energy are the same from RHIC to LHC.

In the left (right) panel of Fig. 3.45 the ratio K^{*0}/K and ϕ/K is plotted as a function of collision energy ($\sqrt{s_{NN}}$) for different collision systems. The K^{*0}/K and ϕ/K ratios in pp collisions are found to be independent of the collision energy and of the colliding system. The only exception is the K^{*0} in Pb–Pb collisions; we attribute the suppression of the K^{*0}/K ratio to final state effects in the late hadronic stage [52]. The behaviours of these ratios in pp collisions agree with the predictions [52, 59] of a thermal model in the grand-canonical limit.

The ϕ/K^{*0} as a function of center-of-mass energy is plotted in Fig. 3.46. Despite K^{*0} and ϕ having similar masses, the $\phi/K^{*0} \neq 1$. This hints at different particle production mechanisms for K^{*0} and ϕ .

3.6.4 Comparison to models

QCD-inspired Monte Carlo (MC) event generators like PYTHIA 8 [36], PHOJET [73, 74] and EPOS-LHC [37] are used to study multi-particle production, which is predominantly a soft, non-perturbative process. Our measurements are compared with these MC model predictions. PYTHIA 8 and PHOJET use the Lund string

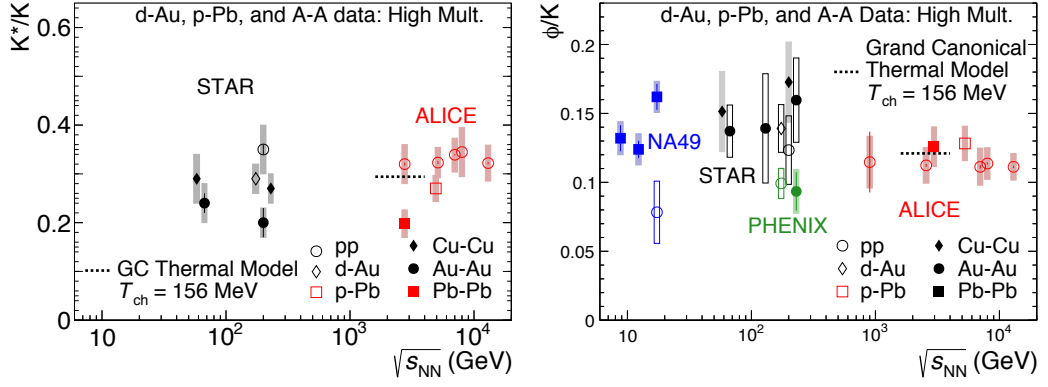


Figure 3.45: (Color online) Particle ratios of K^{*0}/K (left) and ϕ/K (right) are presented for pp, high-multiplicity p-Pb, central d-Au, and central A-A collisions [38–40] as a function of the collision energy. Bars (when present) represent statistical uncertainties. Boxes represent the total systematic uncertainties or the total uncertainties for cases when separate statistical uncertainties were not reported. The value given by a grand-canonical thermal model with a chemical freeze-out temperature of 156 MeV [59] is also shown.

fragmentation model [69] for the hadronisation of light and heavy quarks. We compare our data with the Monash 2013 tune [36] for PYTHIA 8, which uses an updated parameter set for Lund hadronisation compared to previous tunes. To describe the non-perturbative phenomena (soft/semi-hard processes), PYTHIA 8 includes multiple parton-parton interactions while PHOJET uses the Dual Parton Model. For hard scatterings, particle production in both models is based on perturbative QCD. For multiple scatterings, the EPOS-LHC model [37] invokes Gribov’s Reggeon Field Theory [70], which features a collective hadronisation via the core-corona mechanism [71]. The final state partonic system consists of longitudinal flux tubes which fragment into string segments. The high energy density string-segments form so-called “core” region, which evolves hydrodynamically to form the bulk part of the system in the final state. The low-density region is known as the “corona”, which expands and breaks via the production of quark-antiquark pairs and hadronises using the unmodified string fragmentation. Recent data from LHC are used to tune the EPOS-LHC model.

Figure. 3.47 shows a comparison of the K^{*0} (top) and ϕ (bottom) p_T spectra in inelastic pp collisions with PYTHIA8 (Monash tune), PHOJET and EPOS-

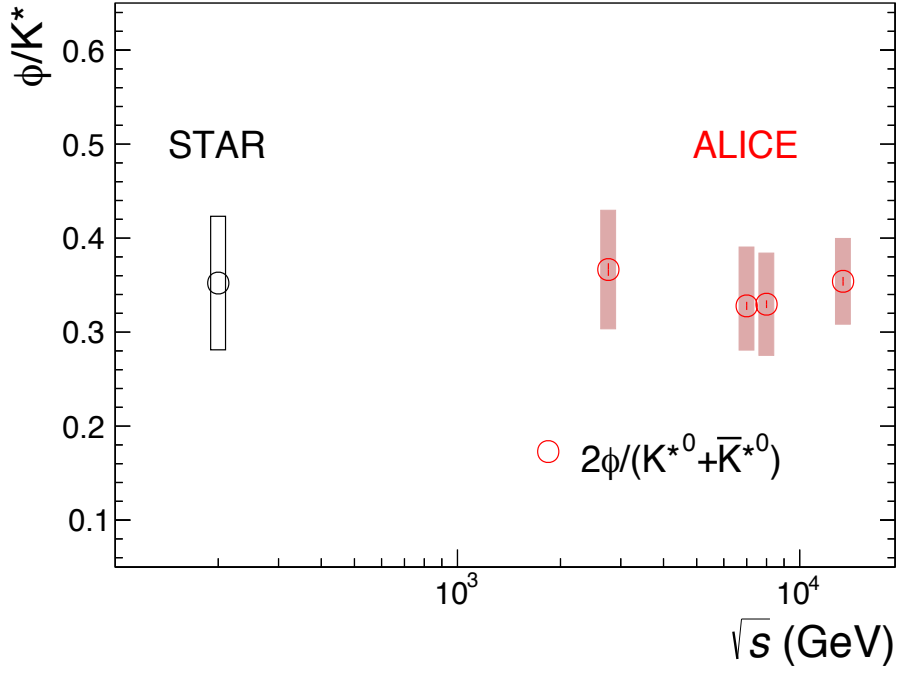


Figure 3.46: (Color online) Particle ratio ϕ/K^{*0} presented for pp collisions [38–40] as a function of the collision energy. Bars (when present) represent statistical uncertainties. Boxes represent the total systematic uncertainties or the total uncertainties for cases when separate statistical uncertainties were not reported.

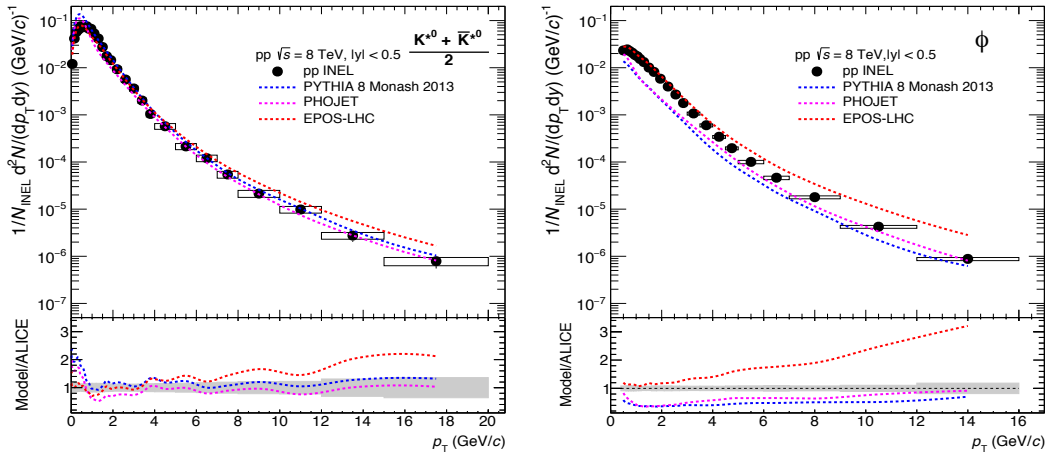


Figure 3.47: (Color online) Comparison of the K^{*0} (left) and ϕ (right) p_T spectra in inelastic pp collisions obtained from PYTHIA8 (Monash2013 tune), PHOJET and EPOS-LHC [38–40]. The bottom plots show the ratios of p_T spectra from the models to the p_T spectra measured by ALICE. The total fractional uncertainties from data are shown as shaded boxes.

LHC. The bottom panels show the ratios of p_T spectra from models to the p_T spectra measured by ALICE. The total fractional uncertainties from the real data, including both statistical and systematic uncertainties are shown as shaded boxes. PYTHIA 8 (Monash2013 tune) describes the p_T spectra for K^{*0} at very low p_T but underestimates it at intermediate p_T and approaches the experimental data at high p_T . For the ϕ meson, PYTHIA 8 (Monash tune) predicts the yields within 40-60 % below the experimental data. PHOJET has a softer p_T spectra for K^{*0} and it explains the data above $p_T > 4$ GeV/ c within the uncertainties. For the ϕ meson, PHOJET predicts the yields similar to PHYTHIA 8 at low p_T , while it approaches the experimental data at higher p_T . For K^{*0} , EPOS-LHC describes the p_T spectra at low p_T and overestimates the data above 4 GeV/ c . For the ϕ meson, the EPOS-LHC model agrees marginally with data at low p_T and deviates monotonically from data with increasing p_T .

3.7 Summary

In this chapter, the measurements of K^{*0} and ϕ production at mid-rapidity in minimum bias pp collisions at $\sqrt{s} = 8$ TeV in the range $0 < p_T < 20$ GeV/ c for K^{*0} and $0.4 < p_T < 16$ GeV/ c for ϕ has been presented. In comparison to other LHC energies, hardening of the p_T spectra is observed with increase in collision energy. The K^{*0}/π and ϕ/π ratios seem to be independent of collision energy within systematic uncertainties. This indicates that there is no strangeness enhancement in minimum bias pp collisions as the collision energy is increased. Similar behavior is observed for the K^{*0}/K and ϕ/K ratios as a function of collision energy. Also, no significant energy dependence of the ϕ/K^{*0} ratio in minimum bias pp collisions at LHC energies is observed, which suggests there is no strong energy dependence of the chemistry of the system. The p-Pb measurements at $\sqrt{s_{NN}} = 8.16$ TeV are in preparation and these pp results will serve as baseline for those in p-Pb collisions.

Chapter 4

Application of Non-Extensive Statistics in High Energy Physics

High multiplicity pp-collisions at LHC give us the opportunity to study matter under extreme conditions *i.e.* at high temperature and/or energy density, which earlier was perceived to be created only in central nuclear collisions. The kinetic freeze-out surface of identified particles can be determined by studying the transverse momentum (p_T) spectra of the produced particles in high energy nuclear and hadronic collisions. Higher probability of multi-partonic interactions at higher collision energies in pp-collisions is believed to be responsible for high-multiplicity events in pp-collisions. These multi-partonic interactions might lead to thermalization in high multiplicity pp-collisions, which could possibly be described by statistical models. The non-exponential behaviour of the transverse momentum spectra at high- p_T is not explained by the Boltzmann-Gibbs distribution function. This hints for a system, which is away from thermal equilibrium. It is now established by various RHIC and LHC experiments that the p_T spectra in pp-collisions is described by the Tsallis non-extensive statistics. We use a thermodynamically consistent Tsallis non-extensive distribution function, which behaves as an exponential function at low- p_T and a power-law function at high- p_T . The parameter q in the non-extensive statistics is a measure of the deviation of the produced system from an equilibrium distribution function. Similarly, we can investigate any possible hint of collectivity in small systems by looking in to the low- p_T part of the transverse spectra using the Boltzmann-Gibbs Blast Wave

(BGBW) model. To explain the identified particle spectra in heavy-ion collisions, incorporation of radial flow in Tsallis statistics is necessary and thus assuming $(q - 1)$ being close to zero, a Taylor series expansion of Tsallis distribution function in powers of $(q - 1)$ is provided. Similarly, the equation of state (EoS) of the system produced in high energy collisions can be understood by looking into the speed of sound in the medium. This also helps in looking for any critical behavior the system could undergo during its space-time evolution. The study of the speed of sound in hadronic matter hence becomes a subject of paramount importance in heavy-ion collisions. Let's first discuss about the Tsallis non-extensive statistics.

4.1 Tsallis Non-extensive Statistics

For a wide range of physical systems, Boltzmann–Gibbs (BG) statistics has been serving as founding structure especially for systems containing a large number of particles, in presence of short range correlations (exponentially decaying). Systems beyond such short range correlation, where the memory effects are significant and non-Markovian ¹ processes are likely to occur require a theoretical modelling. A generalised statistical approach needs to address systems under consideration such that in an appropriate limit the generalised statistics reduces to the BG statistics. Tsallis statistics [48] which is a generalisation of BG statistics has been very successful in describing systems with long range correlation with the generalised entropy

$$S_q = \frac{1 - \sum_{i=1}^W P_i^q}{q - 1}, (q \in \mathbb{R}; \sum_{i=1}^W P_i = 1), \quad (4.1)$$

or, for a continuum distribution

$$S_q = \frac{1 - \text{Tr} \rho^q}{q - 1}. \quad (4.2)$$

Here P_i are the probabilities of micro states, q is a positive integer and ρ is the corresponding density operator. The non-extensivity parameter q tells

¹A non-Markovian process is a stochastic process that doesn't exhibit the Markov property—a Markov process is a random process in which the future is independent of the past, given the present.

about the deviation of thermodynamic systems from an equilibrium distribution. The physical understanding of the non-extensive parameter has become crucial importance in Tsallis statistics and its applications to high-energy physics. The name non-extensivity arises from the non-additive nature of the entropy, which is given by

$$S_q(A + B) = S_q(A) + S_q(B) + (1 - q)S_q(A)S_q(B), \quad (4.3)$$

where A and B are two independent systems in the sense that the probabilities of the microstates of A+B factorize into those of A and B. The additive nature of the entropy can be restored when $q \rightarrow 1$, which is the case for the BG statistics.

A modified form of Tsallis distribution function is given by,

$$f = \left[1 + (q - 1) \frac{E - \mu}{T} \right]^{-\frac{1}{q-1}}, \quad (4.4)$$

where T and μ are the temperature and the chemical potential. In order to achieve thermodynamic consistency an extra power of q is required in the distribution function:

$$f^q = \left[1 + (q - 1) \frac{E - \mu}{T} \right]^{-\frac{q}{q-1}}, \quad (4.5)$$

Numerically the difference between Eq. 4.4 and 4.5 is not large as the value of q is always close to 1 for hadronic and nuclear collisions at relativistic energies. Since the year 2012, this modified form of Tsallis distribution function (Eq. 4.5) [75] has been widely used in the community and is popularly termed as “*thermodynamically consistent Tsallis distribution function*”.

It can be shown explicitly [76] that the expressions for entropy, S , particle number, N , energy density, ϵ , and the pressure, P in the context of this thermodynamically consistent picture of Tsallis non-extensive approach, are given

by,

$$S = -gV \int \frac{d^3p}{(2\pi)^3} [f^q \ln_q f - f], \quad (4.6)$$

$$N = gV \int \frac{d^3p}{(2\pi)^3} f^q, \quad (4.7)$$

$$\epsilon = g \int \frac{d^3p}{(2\pi)^3} E f^q, \quad (4.8)$$

$$P = g \int \frac{d^3p}{(2\pi)^3} \frac{p^2}{3E} f^q, \quad (4.9)$$

where V is the volume and g is the degeneracy factor.

The function appearing in Eq. 4.6 is often referred to as q -logarithm and is defined by

$$\ln_q(x) \equiv \frac{x^{1-q} - 1}{1 - q}.$$

The first and second laws of thermodynamics lead to the following two differential relations:

$$d\epsilon = T ds + \mu dn, \quad (4.10)$$

$$dP = s dT + n d\mu. \quad (4.11)$$

where, $s = S/V$ and $n = N/V$ are the entropy and particle number densities, respectively.

Thermodynamic consistency conditions which require that the following relations to be satisfied need to use f^q instead of f to define the thermodynamic variables.

$$T = \left. \frac{\partial \epsilon}{\partial s} \right|_n, \quad (4.12)$$

$$\mu = \left. \frac{\partial \epsilon}{\partial n} \right|_s, \quad (4.13)$$

$$n = \left. \frac{\partial P}{\partial \mu} \right|_T, \quad (4.14)$$

$$s = \left. \frac{\partial P}{\partial T} \right|_\mu. \quad (4.15)$$

Eq. 4.12, in particular, shows that the variable T appearing in Eq. 4.4 can indeed be identified as a thermodynamic temperature and is more than just

another parameter. It is straightforward to show that these relations are indeed satisfied [76].

The identified particle distribution can be rewritten, using variables appropriate for high-energy physics based on the above expressions as

$$\frac{d^2N}{p_T dp_T dy} = \frac{gV}{(2\pi)^2} m_T \cosh y \left[1 + (q-1) \frac{m_T \cosh y - \mu}{T} \right]^{-\frac{q}{q-1}} \quad (4.16)$$

At mid-rapidity, $y = 0$, and LHC energies where $\mu \sim 0$, the transverse momentum distribution is given by

$$\frac{d^2N}{p_T dp_T dy} = \frac{gV}{(2\pi)^2} m_T \left[1 + (q-1) \frac{m_T}{T} \right]^{-\frac{q}{q-1}}. \quad (4.17)$$

We have used eq. 4.17 to analyse the transverse momentum spectra of the identified particles spectra for different multiplicity classes in pp collisions at $\sqrt{s} = 7$ TeV.

4.2 Transverse Momentum Spectra of Identified Hadrons using Tsallis Non-extensive Statistics

In this work [77], we analyse the transverse momentum spectra of $\pi^\pm$, $K^\pm$, K_S^0 , $p + \bar{p}$, $K^{*0} + \bar{K}^{*0}$, ϕ , $\Lambda + \bar{\Lambda}$, $\Xi^- + \bar{\Xi}^+$ and $\Omega^- + \bar{\Omega}^+$ produced in pp-collisions at $\sqrt{s} = 7$ TeV at the LHC, measured by the ALICE [78, 79].

We use Eq. 4.17 to fit the transverse momentum spectra of various particles measured experimentally for different multiplicity classes at $\sqrt{s} = 7$ TeV. The fittings are shown in Figs. 4.1 and 4.2. The fitting is performed using the TMinuit class available in ROOT library keeping all the parameters free. Here, we follow the notion of a mass dependent differential freeze-out scenario [80, 81], where particles freeze-out at different times, which correspond to different system volumes and temperatures. Henceforth, we study the thermodynamic parameters in the context of non-extensive statistics. After fitting we found that, χ^2/ndf is below 1 for all the considered particles for highest multiplicity while it increases

as the multiplicity decreases. This shows that the spectra are very well described by the non-extensive statistics particularly at the highest multiplicity class.

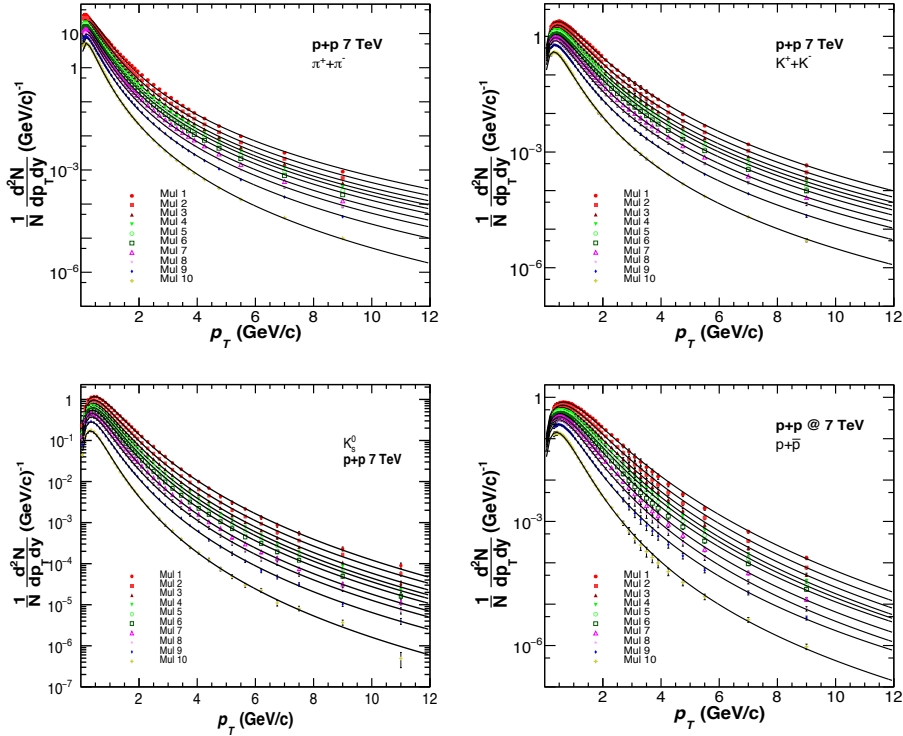


Figure 4.1: (Color online) Experimentally measured p_T -spectra of $\pi^\pm$, $K^\pm$, K_S^0 and $(p + \bar{p})$ [78], fitted with non-extensive Tsallis distribution for pp-collisions at $\sqrt{s} = 7$ TeV using Eq. 4.17 for various multiplicity classes [77]. The extracted parameters are given in table 4.2.

Figure 4.3 represents the temperature parameter, T extracted in the fitting as a function of event multiplicity for all the considered particles. We notice a monotonic increase in T with the increase in particle multiplicity for all the hadrons. For heavier particles, the temperature is observed to be higher, which indicates an early freeze-out of the these particles. We also find that the temperature for lighter particles does not change appreciably with multiplicity but for heavier particles it shows a significant variation of T with charged particle multiplicity.

We have also shown the variation of the non-extensive parameter, q with charged particle multiplicity in Fig. 4.4. The value of q decreases monotonically for higher multiplicity classes for all the particles discussed in this work. Theses

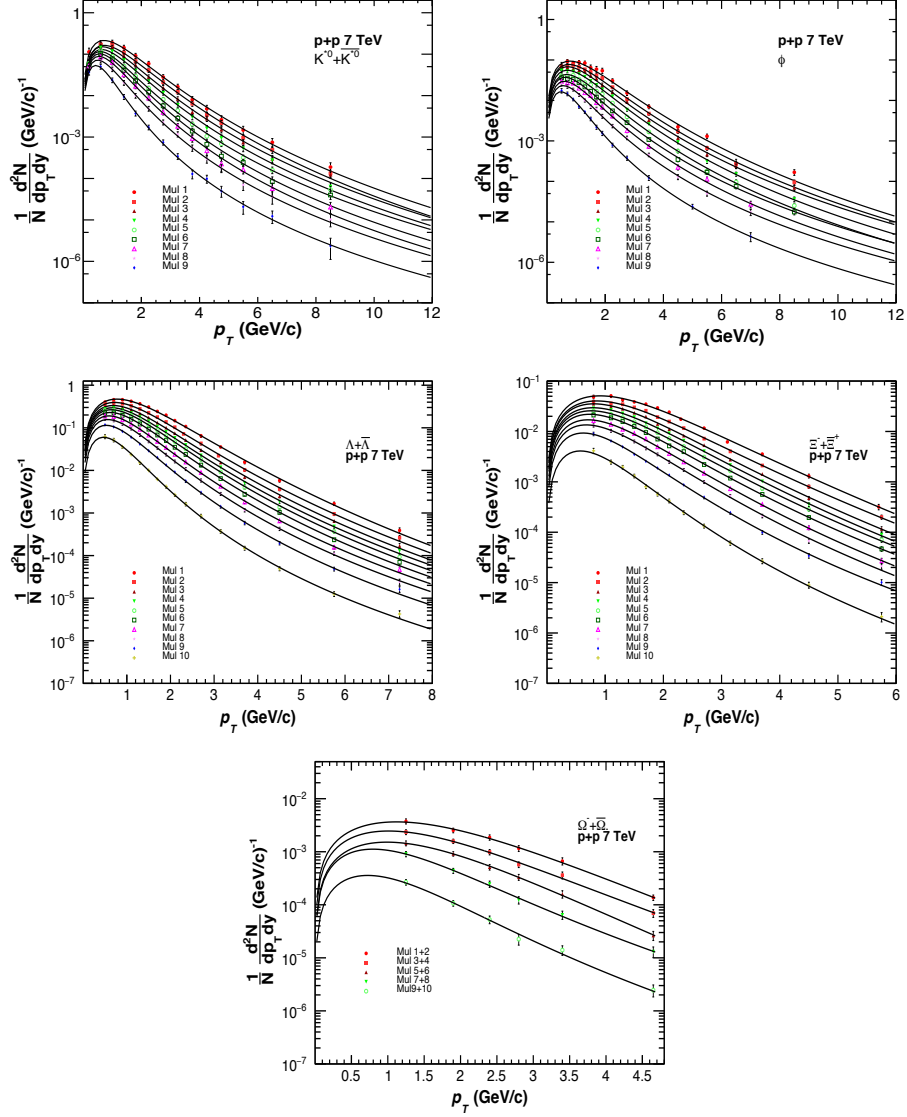


Figure 4.2: (Color online) The p_T -spectra of $K^{*0} + \bar{K}^{*0}$, ϕ , $\Lambda + \bar{\Lambda}$ [78], $\Xi^- + \bar{\Xi}^+$ and $\Omega^- + \bar{\Omega}^+$ fitted with Tsallis distribution for pp-collisions at $\sqrt{s} = 7$ TeV using Eq. 4.17 for various multiplicity classes [77]. The extracted parameters are given in table 4.2.

findings suggest that the system formed in higher multiplicity class is close to thermal equilibrium. However for $\pi^\pm$, $K^\pm$ and K_S^0 , q is almost independent of the charged particle multiplicity density. The fact that the q -values go on decreasing with multiplicity is an indicative of the tendency of the produced systems towards thermodynamic equilibrium. This goes inline with the naive expectations while understanding the microscopic view of systems approaching thermodynamic equilibrium. A similar tendency of q decreasing with number

Table 4.1: The extracted Tsallis parameters as well as the χ^2/ndf for all the multiplicity classes. For $\Omega^- + \bar{\Omega}^+$, multiplicity classes are combined to deal with the statistics [79], For K_S^0 , $\Lambda + \bar{\Lambda}$, $\Xi^- + \bar{\Xi}^+$ and $\Omega^- + \bar{\Omega}^+$, the parameters are taken from our earlier work [77].

Particles		Multiplicity class									
		Mul1	Mul2	Mul3	Mul4	Mul5	Mul6	Mul7	Mul8	Mul9	Mul10
$\pi^\pm$	T (GeV)	0.093 ± 0.001	0.089 ± 0.001	0.087 ± 0.001	0.085 ± 0.001	0.083 ± 0.001	0.081 ± 0.001	0.078 ± 0.001	0.076 ± 0.001	0.073 ± 0.001	0.068 ± 0.001
	q	1.163 ± 0.001	1.162 ± 0.001	1.161 ± 0.001	1.161 ± 0.001	1.160 ± 0.001	1.159 ± 0.001	1.157 ± 0.001	1.156 ± 0.001	1.152 ± 0.001	1.143 ± 0.001
	χ^2/ndf	7.574	7.757	6.792	5.950	5.294	4.349	3.074	1.929	0.603	0.663
$K^\pm$	T (GeV)	0.148 ± 0.003	0.135 ± 0.003	0.125 ± 0.003	0.118 ± 0.003	0.113 ± 0.003	0.107 ± 0.003	0.098 ± 0.003	0.090 ± 0.003	0.080 ± 0.002	0.060 ± 0.002
	q	1.143 ± 0.002	1.146 ± 0.002	1.148 ± 0.002	1.149 ± 0.002	1.150 ± 0.002	1.151 ± 0.002	1.152 ± 0.002	1.153 ± 0.002	1.153 ± 0.002	1.150 ± 0.002
	χ^2/ndf	0.290	0.236	0.255	0.200	0.155	0.130	0.131	0.084	0.062	0.145
$p + \bar{p}$	T (GeV)	0.183 ± 0.006	0.160 ± 0.005	0.143 ± 0.005	0.131 ± 0.005	0.119 ± 0.005	0.101 ± 0.004	0.090 ± 0.004	0.075 ± 0.004	0.049 ± 0.004	0.017 ± 0.001
	q	1.112 ± 0.003	1.115 ± 0.002	1.117 ± 0.002	1.118 ± 0.002	1.121 ± 0.002	1.126 ± 0.002	1.126 ± 0.002	1.127 ± 0.002	1.135 ± 0.002	1.136 ± 0.001
	χ^2/ndf	0.483	0.701	0.622	0.384	0.554	0.637	0.452	0.297	0.501	0.331
K_S^0	T (GeV)	0.152 ± 0.001	0.137 ± 0.001	0.131 ± 0.004	0.124 ± 0.003	0.119 ± 0.005	0.111 ± 0.004	0.103 ± 0.004	0.095 ± 0.002	0.085 ± 0.003	0.068 ± 0.003
	q	1.141 ± 0.001	1.144 ± 0.001	1.144 ± 0.002	1.145 ± 0.002	1.146 ± 0.003	1.148 ± 0.002	1.148 ± 0.002	1.150 ± 0.002	1.150 ± 0.002	1.147 ± 0.002
	χ^2/ndf	0.275	0.429	0.235	0.323	0.309	0.286	0.484	0.384	0.321	0.494
$\Lambda + \bar{\Lambda}$	T (GeV)	0.245 ± 0.0	0.201 ± 0.0	0.179 ± 0.0	0.159 ± 0.0	0.146 ± 0.0	0.128 ± 0.0	0.102 ± 0.0	0.082 ± 0.0	0.056 ± 0.0	0.010 ± 0.0
	q	1.086 ± 0.006	1.097 ± 0.004	1.101 ± 0.007	1.106 ± 0.004	1.108 ± 0.004	1.111 ± 0.001	1.118 ± 0.003	1.123 ± 0.002	1.128 ± 0.004	1.139 ± 0.001
	χ^2/ndf	0.554	0.543	0.281	0.311	0.272	0.307	0.201	0.160	0.312	0.248
$\Xi^- + \bar{\Xi}^+$	T (GeV)	0.308 ± 0.0	0.260 ± 0.0	0.224 ± 0.0	0.212 ± 0.0	0.186 ± 0.0	0.164 ± 0.0	0.147 ± 0.0	0.122 ± 0.0	0.074 ± 0.0	0.045 ± 0.001
	q	1.069 ± 0.015	1.081 ± 0.005	1.086 ± 0.005	1.088 ± 0.004	1.096 ± 0.003	1.100 ± 0.003	1.101 ± 0.003	1.108 ± 0.002	1.121 ± 0.002	1.122 ± 0.002
	χ^2/ndf	0.837	0.458	0.350	0.133	0.168	0.232	0.237	0.543	0.313	0.369
		Mul1	Mul2	Mul3	Mul[4 + 5]		Mul6	Mul7	Mul8	Mul9	Mul10
Φ	T (GeV)	0.241 ± 0.027	0.212 ± 0.018	0.181 ± 0.017	0.177 ± 0.015		0.161 ± 0.015	0.116 ± 0.015	0.109 ± 0.018	0.070 ± 0.018	0.019 ± 0.002
	q	1.116 ± 0.012	1.122 ± 0.008	1.131 ± 0.008	1.127 ± 0.006		1.128 ± 0.007	1.145 ± 0.007	1.141 ± 0.009	1.153 ± 0.009	1.159 ± 0.002
	χ^2/ndf	1.790	0.687	0.840	0.268		0.666	0.293	0.479	0.479	0.159
$K^{*0} + \bar{K}^{*0}$	T (GeV)	0.225 ± 0.029	0.230 ± 0.021	0.185 ± 0.020	0.172 ± 0.018		0.132 ± 0.017	0.124 ± 0.019	0.109 ± 0.019	0.076 ± 0.017	0.024 ± 0.019
	q	1.117 ± 0.013	1.112 ± 0.009	1.126 ± 0.009	1.126 ± 0.009		1.137 ± 0.008	1.134 ± 0.010	1.135 ± 0.010	1.145 ± 0.009	1.155 ± 0.011
	χ^2/ndf	0.394	0.615	0.806	0.681		0.587	0.807	0.339	0.165	0.268
		Mul[1+2]		Mul[3+4]		Mul[5+6]		Mul[7+8]		Mul[9+10]	
$\Omega^- + \bar{\Omega}^+$	T (GeV)	0.422 ± 0.0		0.310 ± 0.0		0.330 ± 0.001		0.104 ± 0.0		0.052 ± 0.0	
	q	1.035 ± 0.028		1.066 ± 0.011		1.044 ± 0.005		1.117 ± 0.015		1.124 ± 0.006	
	χ^2/ndf	0.297		0.478		0.092		0.238		0.652	

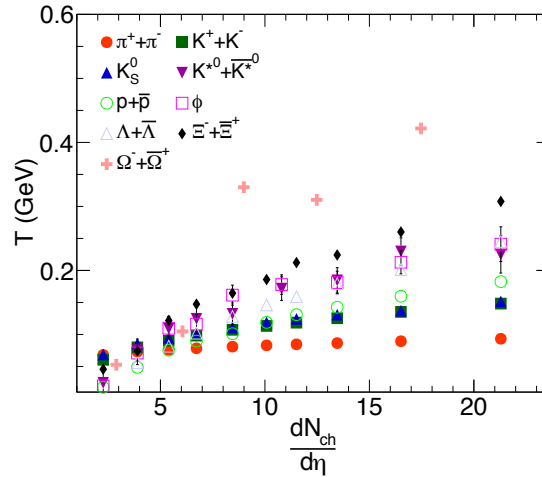


Figure 4.3: (Color online) Multiplicity dependence of T for pp-collisions at $\sqrt{s} = 7$ TeV using Eq. 4.17 as a fitting function [77].

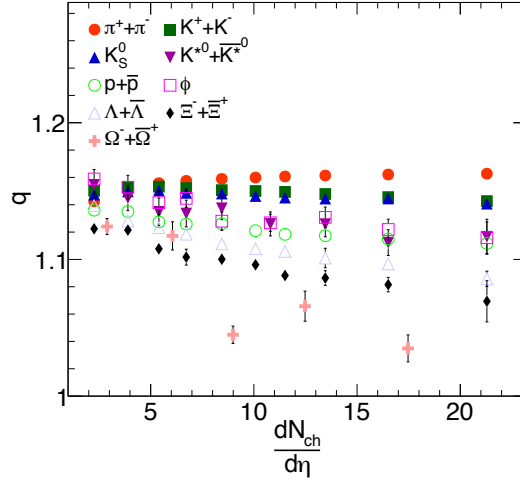


Figure 4.4: (Color online) Multiplicity dependence of the non-extensive parameter, q for pp-collisions at $\sqrt{s} = 7$ TeV using Eq. 4.17 as a fitting function [77].

of participating nucleons for Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV has been observed for the bulk part ($p_T < 6$ GeV/ c) of the charged hadron spectra [82, 83]. The present study is useful in understanding the microscopic features of degrees of equilibration and their dependencies on the number of particles in the system.

It has been manifested in this work, the Tsallis distribution provides a complete description of identified particle spectra produced in pp-collisions at $\sqrt{s} = 7$ TeV upto very high- p_T . The variable T shows a systematic increase with multiplicity, the heaviest baryons showing the steepest increase. This is an indication of a mass hierarchy in particle freeze-out. The obtained parameters show variations with the event multiplicity. The notable variation of the non-extensive parameter, q which decreases towards the value 1 as the multiplicity increases and this effect is more significant for heavy mass particles. This shows the tendency of the produced system to equilibrate with higher multiplicities. This goes inline with the expected multi-partonic interactions, which increase for higher multiplicities in pp-collisions and are thus responsible for bringing the system towards thermodynamic equilibrium [84]. Here, we want to mention that the extracted parameters for ω have more fluctuations and this might be due to the lower statistics. Recent results on high-multiplicity pp-collisions such as collectivity in small systems, and enhancement of yields of strange and multi-strange particles in com-

parison to pion as a function of charge particle multiplicity, which is similar to what we find in heavy-ion collisions opens up new challenges and this needs to be understood. Thus, we can investigate any possible hint of collectivity by looking into the low- p_T part of the transverse spectra using the Boltzmann-Gibbs Blast Wave (BGBW) model.

4.3 Transverse Momentum Spectra of Identified Hadrons Using Boltzmann-Gibbs Blast Wave (BGBW) Model

The expression for invariant yield in the framework of BGBW is given as follows [85]:

$$E \frac{d^3 N}{dp^3} = D \int d^3 \sigma_\mu p^\mu \exp\left(-\frac{p^\mu u_\mu}{T}\right), \quad (4.18)$$

where the particle four-momentum is,

$$p^\mu = (m_T \cosh y, p_T \cos \phi, p_T \sin \phi, m_T \sinh y), \quad (4.19)$$

the four-velocity is given by,

$$u^\mu = \cosh \rho (\cosh \eta, \tanh \rho \cos \phi_r, \tanh \rho \sin \phi_r, \sinh \eta), \quad (4.20)$$

while the kinetic freeze-out surface is parametrised as,

$$d^3 \sigma_\mu = (\cosh \eta, 0, 0, -\sinh \eta) \tau r dr d\eta d\phi_r. \quad (4.21)$$

Here, η is the space-time rapidity. With simplification assuming Bjorken correlation in rapidity, *i.e.* $y = \eta$ [86], Eq. 4.18 can be expressed as:

$$\left. \frac{d^2 N}{dp_T dy} \right|_{y=0} = D \int_0^{R_0} r dr K_1\left(\frac{m_T \cosh \rho}{T_{\text{kin}}}\right) I_0\left(\frac{p_T \sinh \rho}{T_{\text{kin}}}\right), \quad (4.22)$$

where D is the normalisation constant. Here g is the degeneracy factor and $m_T = \sqrt{p_T^2 + m^2}$ is the transverse mass. $K_1\left(\frac{m_T \cosh \rho}{T_{\text{kin}}}\right)$ and $I_0\left(\frac{p_T \sinh \rho}{T_{\text{kin}}}\right)$ are the modified Bessel's functions and are given by,

$$K_1\left(\frac{m_T \cosh \rho}{T}\right) = \int_0^\infty \cosh y \exp\left(-\frac{m_T \cosh y \cosh \rho}{T_{\text{kin}}}\right) dy,$$

$$I_0\left(\frac{p_T \sinh \rho}{T}\right) = \frac{1}{2\pi} \int_0^{2\pi} \exp\left(\frac{p_T \sinh \rho \cos \phi}{T_{\text{kin}}}\right) d\phi,$$

where ρ in the integrand is a parameter given by $\rho = \tanh^{-1} \beta$, with $\beta = \beta_s \left(\xi\right)^n$ [85, 87–89] is the radial flow. β_s is the maximum surface velocity and $\xi = \left(r/R_0\right)$, with r as the radial distance. In the blast-wave model the particles closer to the centre of the fireball move slower than the ones at the edges. The average of the transverse velocity can be evaluated as [90],

$$\langle \beta \rangle = \frac{\int \beta_s \xi^n \xi d\xi}{\int \xi d\xi} = \left(\frac{2}{2+n}\right) \beta_s. \quad (4.23)$$

In our calculation, we use a linear velocity profile, ($n = 1$) and R_0 is the maximum radius of the expanding source at freeze-out ($0 < \xi < 1$).

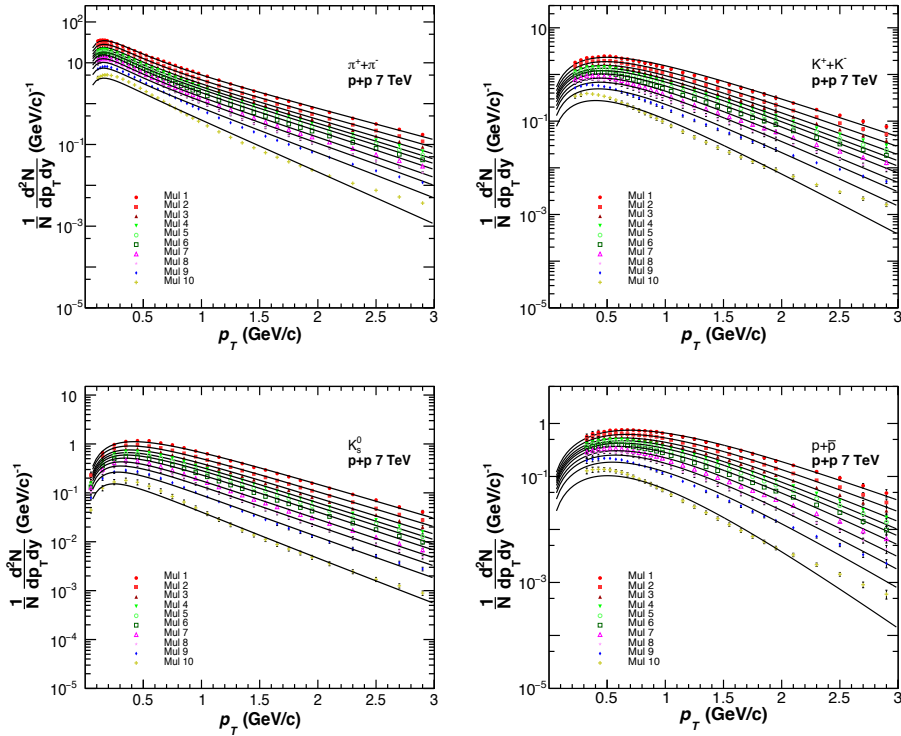


Figure 4.5: (Color online) Fitting of experimentally measured p_T -spectra of $\pi^\pm$, $K^\pm, K_S^0$ and $p + \bar{p}$ with BGBW model for pp-collisions at $\sqrt{s} = 7$ TeV using Eq. 4.22 for various multiplicity classes [91].

In figure 4.5 and 4.6, we show the fitting of p_T -spectra of $\pi^\pm$, $K^\pm$, K_S^0 , $p + \bar{p}$, $K^{*0} + \bar{K}^{*0}$, ϕ and $\Lambda + \bar{\Lambda}$ measured in pp-collisions at $\sqrt{s} = 7$ TeV in

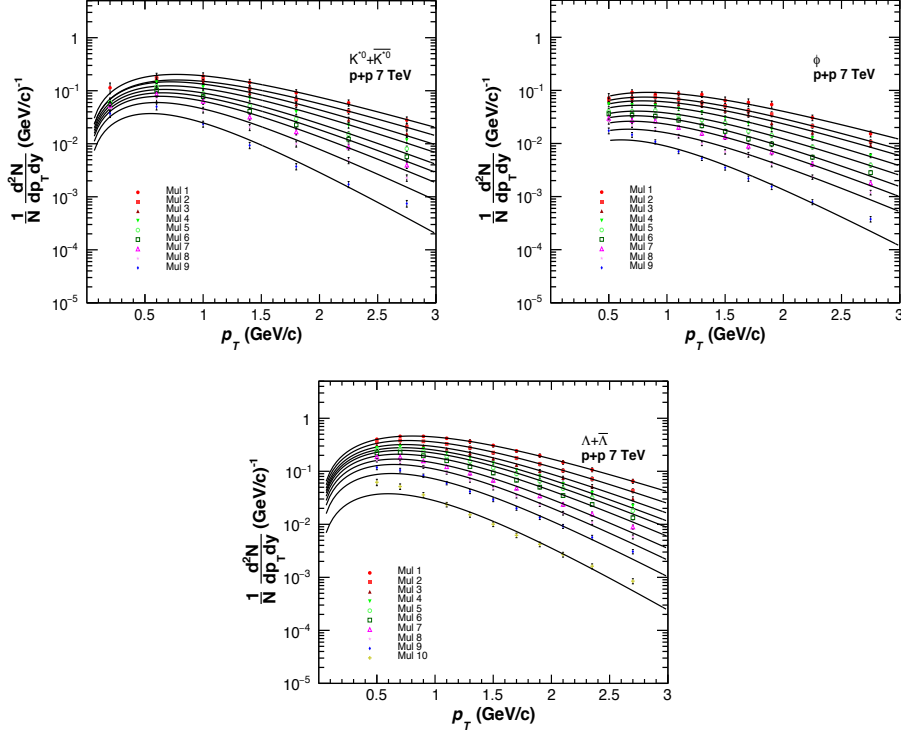


Figure 4.6: (Color online) Fitting of experimentally measured p_T -spectra of $K^{*0} + \bar{K}^{*0}$, ϕ and $\Lambda + \bar{\Lambda}$ with BGBW model for pp-collisions at $\sqrt{s} = 7$ TeV using Eq. 4.22 for various multiplicity classes [91].

various multiplicity classes. Here, we have used BGBW function to fit the spectra given by Eq. 4.22, where T_{kin} and radial flow (β) are the free parameters. We perceive that BGBW function explains the experimental data very well for all the particles at lower p_T (~ 3 GeV/c) with a good χ^2/ndf . However, a higher χ^2/ndf is obtained in case of $\pi^\pm$ due to the underestimation of experimental data at lower p_T , as we are not considering the contributions from resonances.

We have extracted the kinetic freeze-out temperature and radial flow velocity for all the hadrons considered here. Figure 4.7 represents the extracted T_{kin} as a function of charged particle multiplicity for identified hadrons. We notice that, T_{kin} is higher for massive particles in comparison to lighter ones which supports the differential freeze out scenario and suggests that massive particles freeze-out earlier from the system. However, K_S^0 behaves differently in the system. It is also evident from the figure that the particles having similar masses are coming from the same freeze-out hypersurface at highest multiplicity class.

Further, we observe multiplicity independent behaviour of T_{kin} particularly for lighter hadrons.

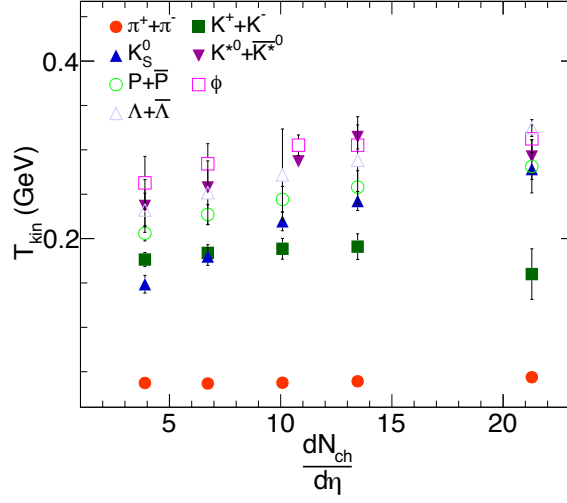


Figure 4.7: (color online) Multiplicity dependence of T_{kin} for pp-collisions at $\sqrt{s} = 7$ TeV using Eq. 4.22 as the fitting function [91].

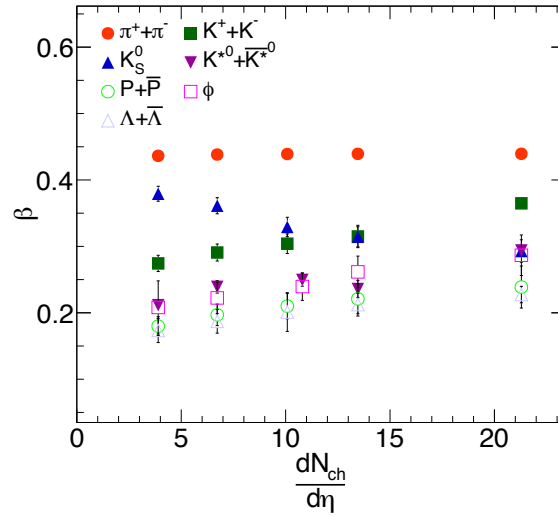


Figure 4.8: (Color online) Multiplicity dependence of the radial flow, β for pp-collisions at $\sqrt{s} = 7$ TeV using Eq. 4.22 as the fitting function [91].

In figure 4.8, we have demonstrated the extracted radial flow velocity for the identified particles as a function of multiplicity class. We find that β value reduces as the particle mass increases in all the multiplicity classes except for

proton. Again, K_S^0 behaves in a different manner which is not understood in the present work. Further, almost equal magnitude of radial flow is observed for K^{*0} and ϕ , which have similar masses. Furthermore, we observe that p has mass similar to K^{*0} and ϕ but the radial flow of p is closer to Λ , as both are baryons, which advocates that baryons like p and Λ freeze-out from the same hypersurface while mesons such as ϕ and K^{*0} freeze-out from a different hypersurface.

The observed radial flow velocity in a differential freeze-out scenario in pp collisions at $\sqrt{s} = 7$ TeV, on the average is found to be around 0.4. For a comparison with LHC $\sqrt{s_{NN}} = 2.76$ TeV Pb–Pb central collisions, one observes $\langle\beta\rangle = 0.64$ [92], which is much higher than the collective flow in pp-collisions.

It is experimentally seen that even in heavy-ion collisions, the p_T spectra of identified particles show a deviation from BG-statistics; for which one uses Levy-Tsallis distribution to study particle spectra. In this work, we take thermodynamically consistent Tsallis distribution function and incorporate radial flow to describe the particle spectra in heavy-ion collisions. It is found that the non-extensive parameter, q in nuclear/hadronic collisions is found to be $1 < q < 1.2$ [93]. Assuming $(q - 1)$ close to zero, a Taylor series expansion of Tsallis distribution function in powers of $(q - 1)$ is provided.

4.4 Radial Flow in Non-Extensive Thermodynamics and Study of Particle Spectra at LHC in the Limit of Small $(q - 1)$

4.4.1 Momentum Distributions to First and Second Order in $(q - 1)$

Assuming the parameter q to be close to 1, as is the value for many cases in high energy physics, the modified Tsallis distribution, f^q , appearing in the expressions for the thermodynamic quantities can be expanded in a Taylor series in $q - 1$ with

the following result (see appendix 6.1 for a detailed derivation):

$$\begin{aligned}
& \left[1 + (q-1) \frac{E-\mu}{T} \right]^{-\frac{q}{q-1}} \\
& \simeq e^{-\frac{E-\mu}{T}} \left\{ 1 + (q-1) \frac{1}{2} \frac{E-\mu}{T} \left(-2 + \frac{E-\mu}{T} \right) \right. \\
& + \frac{(q-1)^2}{2!} \frac{1}{12} \left[\frac{E-\mu}{T} \right]^2 \left[24 - 20 \frac{E-\mu}{T} + 3 \left(\frac{E-\mu}{T} \right)^2 \right] \\
& + \mathcal{O} \{ (q-1)^3 \} \\
& + \dots \}
\end{aligned} \tag{4.24}$$

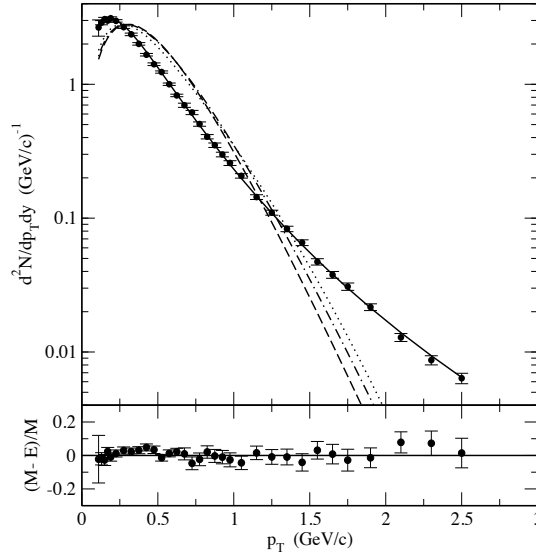


Figure 4.9: Fits to the normalized differential yields of π^+ as measured by the ALICE collaboration in pp-collisions at $\sqrt{s} = 0.9$ TeV [94] fitted with the Tsallis (solid line) and Boltzmann distributions (dashed line). Also shown are fits with the Tsallis distribution keeping terms to first (dash-dotted line) and second order in $(q-1)$ (dotted line). The lower part of the figure shows the difference between model (M) and experiment (E) normalized to the model (M) values [95].

This result can be used for the invariant yield of particles if it is written in terms of the Tsallis distribution,

$$E \frac{dN}{d^3p} = CE \left[1 + (q-1) \frac{E-\mu}{T} \right]^{-\frac{q}{q-1}} \tag{4.25}$$

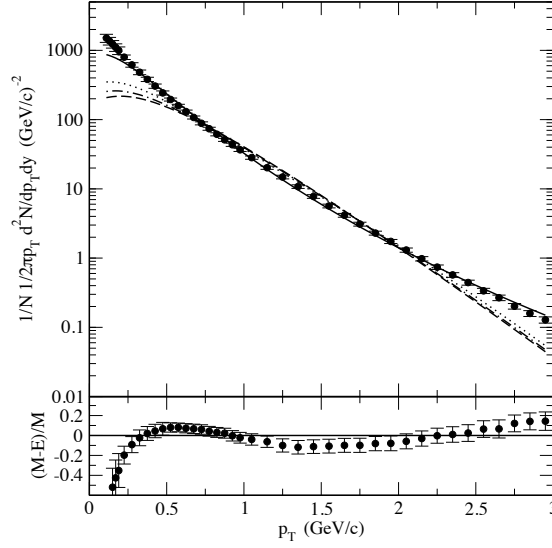


Figure 4.10: Fits to the normalized differential π^- yields as measured by the ALICE collaboration in (0–5)% Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [96] fitted with the Tsallis (solid line) and Boltzmann distributions (dashed line). Also shown are fits with the Tsallis distribution keeping terms to first (dash-dotted line) and second order in $(q-1)$ (dotted line). The lower part of the figure shows the difference between model (M) and experiment (E) normalized to the model (M) values [95].

where $\mathcal{C} \equiv gV/(2\pi)^3$. Let us use the following notations: $(q-1) \equiv x$; $(E-\mu)/T = \Phi$ and $1 + (q-1)\frac{E-\mu}{T} = 1 + x\Phi = f(x)$. Hence, the expansion of the Tsallis distribution up to $\mathcal{O}(x^2)$ can be written as:

$$\begin{aligned} E \frac{dN}{d^3p} &\simeq \mathcal{C} E e^{-\Phi} + \mathcal{C} E \frac{x}{1!} \frac{\Phi}{2} (-2 + \Phi) e^{-\Phi} \\ &+ \mathcal{C} E \frac{x^2}{2!} \frac{\Phi^2}{12} (24 - 20\Phi + 3\Phi^2) e^{-\Phi} \end{aligned} \quad (4.26)$$

Hence one obtains

$$\begin{aligned} \frac{dN}{p_T dp_T dy d\phi} &\simeq \mathcal{C} E e^{-\Phi} + \mathcal{C} E \frac{x}{1!} \frac{\Phi}{2} (-2 + \Phi) e^{-\Phi} \\ &+ \mathcal{C} E \frac{x^2}{2!} \frac{\Phi^2}{12} (24 - 20\Phi + 3\Phi^2) e^{-\Phi} \end{aligned} \quad (4.27)$$

Since, we use a modified form of Tsallis distribution, a comparison with similar work will be worthwhile at this point. With this aim, we integrate over

the rapidity variable to compare the transverse mass spectrum obtained from the present approach with that obtained in Ref. [98]:

$$\begin{aligned}
\frac{dN}{m_T dm_T} = & \frac{gV}{2\pi^2} m_T \left[K_1 \left(\frac{m_T}{T} \right) e^{\frac{\mu}{T}} \right. \\
& - \frac{q-1}{2} \frac{m_T}{T} \left\{ K_0 \left(\frac{m_T}{T} \right) + K_2 \left(\frac{m_T}{T} \right) \right\} \\
& + \frac{q-1}{8} \left(\frac{m_T}{T} \right)^2 \left\{ K_3 \left(\frac{m_T}{T} \right) + 3K_1 \left(\frac{m_T}{T} \right) \right\} \\
& + \frac{\mu(q-1)}{T} K_1 \left(\frac{m_T}{T} \right) + \frac{\mu^2(q-1)}{2T^2} K_1 \left(\frac{m_T}{T} \right) \\
& \left. - \frac{\mu m_T(q-1)}{2T^2} \left\{ K_0 \left(\frac{m_T}{T} \right) + K_2 \left(\frac{m_T}{T} \right) \right\} \right] \\
& + \mathcal{O}((q-1)^2) + \dots
\end{aligned} \tag{4.28}$$

where K_n s are the modified Bessel's functions of second kind (see Appendix 6.4). In principle, while examining the transverse spectra, the rapidity integration should be within a maximum value y_{max} , say. But, owing to the presence of the term $e^{-m_T \cosh y/T}$, the integrand drops down very fast with increasing y . And so, according to the standard practice, we can effectively replace the rapidity integration from 0 to y_{max} to 0 to ∞ so that the integration yields Bessel functions. The first term in the above expression is the well-known formula for a thermal source with a Boltzmann distribution:

$$\frac{dN}{p_T dp_T} = \frac{gV}{2\pi^2} m_T e^{\frac{\mu}{T}} K_1 \left(\frac{m_T}{T} \right) \tag{4.29}$$

In the limit $\mu = 0$, the transverse mass spectrum obtained from Ref. [98] (up to $\mathcal{O}(q-1)$) does not contain the term involving K_0 and K_2 .

The Boltzmann distribution, the pure Tsallis distribution (Eq. 4.17) and the expansion of the Tsallis distribution up to the first and second order of $(q-1)$ (Eq. 4.27) were used to fit the transverse momentum distributions obtained by the ALICE collaboration. The results are shown in Fig. 4.9 for pp collisions at $\sqrt{s} = 0.9$ TeV and in Fig. 4.10 for Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. It is well known that the Tsallis fits give excellent results for pp collisions but are not very good for Pb-Pb collisions as can be seen clearly in Fig. 4.10, where a large deviation at small p_T is seen.

It can be seen from the fits to the pp distribution that the successive terms in $(q-1)$ improve the fits but not in a convincing manner. Clearly, it is the best

to use the full Tsallis distribution and not the series expansion. It might turn out of course that the series expansion could be of use in a different situation where the comparison with a Boltzmann distribution is more relevant. The fits in the figures were done using the MINUIT package with the following numerical results.

In Fig. 4.9 we show fits to the transverse momentum distribution of π^+ in pp collisions at 900 GeV. For the plain Tsallis distribution (solid line) the parameters were obtained as being $T = 70.8$ MeV, $q = 1.1474$. The volume V was determined as corresponding to a spherical radius of 4.81 fm. For the Boltzmann distribution (dashed line) the parameters were determined as being $T = 150.2$ MeV, while the radius used to determine the volume was fixed at a value of 2.65 fm. For the fit using the Boltzmann distribution and the first order term in $(q - 1)$ (dashed-dotted line) the values are $T = 138.4$ MeV, $q = 1.035$ while the radius is given by 2.80 fm. In the last case corresponding to Boltzmann plus first and second orders in $(q - 1)$ (dotted line) one has $T = 121.2$ MeV, $q = 1.065$ and a radius of 3.09 fm. As is well-known and evident, the fit using the Tsallis distribution is very good.

In Fig. 4.10 we show fits to the normalized differential π^- yields in (0–5)% Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV as measured by the ALICE collaboration [96] with the Tsallis (solid line) and Boltzmann distributions (dashed line). Also shown are fits with the Tsallis distribution keeping terms to first order (dash-dotted line) and second order in $(q - 1)$ (dotted line). The lower part of the figure shows the difference between the Tsallis distribution (M) and experiment (E). It is clear that the best fit is achieved with the full Tsallis distribution, whereas, using the Boltzmann distribution the description is not good. Successive corrections in $(q - 1)$ improve the description. There is a clear deviation at very low transverse momentum (below 0.5 GeV) and also at higher values above 2.75 GeV.

4.4.2 Thermodynamic Quantities to First Order in $(q - 1)$

The particle density in Tsallis thermodynamics is given to first order in $(q - 1)$ by the following expression:

$$n^B + (q - 1)n^1 \tag{4.30}$$

where n^B is the standard Boltzmann result for the particle density:

$$n^B = \frac{g}{2\pi^2} e^{\frac{\mu}{T}} T^3 a^2 K_2(a), \quad (4.31)$$

with $a \equiv m/T$, and the first order expression in $q - 1$ is given by

$$\begin{aligned} n^1 = & \frac{ge^{\frac{\mu}{T}} T^3}{4\pi^2} \left[-6a^2 K_2(a) - 2a^3 K_1(a) \right. \\ & - 4a^2 b K_2(a) + 3a^3 K_3(a) + a^4 K_2(a) + a^2 b^2 K_2(a) \\ & \left. - 2a^3 b K_1(a) \right]. \end{aligned} \quad (4.32)$$

In Fig. 4.11 we show the ratio of the particle density to first order in $(q - 1)$ to the full particle density as given by the Tsallis distribution, $(n^B + (q - 1)n^1)/n$ for several values of q indicated in the figure as a function of the temperature T . It can be seen that the expansion in $(q - 1)$ is excellent if $(q - 1) = 0.01$ but rapidly deviates from the full Tsallis distribution for larger values of q . Already for $(q - 1) \approx 0.1$ the deviations are of the order of 10 % as can be seen from Fig. 4.11.

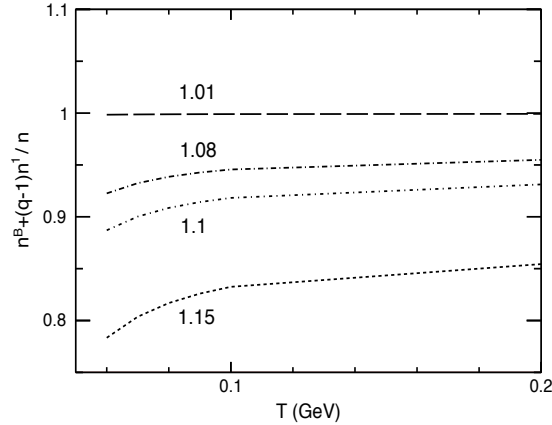


Figure 4.11: The ratio of the particle density calculated to first order in $(q - 1)$ normalized to the particle density of a Tsallis gas as a function of the temperature for different values of the parameter q . The mass is taken as being the pion mass. The values of the parameter q are 1.01 for the dashed line, 1.08 for the dot-(long)dashed line, 1.1 for the dot-(short)dashed line and 1.15 for the dotted line [95].

For comparison we show in Fig. 4.12, the first order expansion compared to the Boltzmann expression, $(n^B + (q - 1)n^1)/n^B$, again as a function of the

temperature T for several values of the parameter q . In this case the deviations are most pronounced for small values of the temperature.

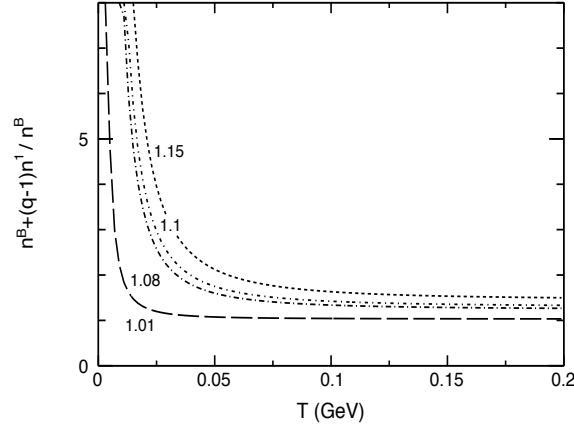


Figure 4.12: The ratio of the particle density calculated to first order in $(q - 1)$ normalized to the particle density of a Boltzmann gas as a function of the temperature for different values of the parameter q . The mass is taken as being the pion mass. The values of the parameter q are 1.01 for the dashed line, 1.08 for the dot-(long)dashed line, 1.1 for the dot-(short)dashed line and 1.15 for the dotted line [95].

Correspondingly, the energy density is obtained as:

$$\epsilon^B + (q - 1)\epsilon^1 \quad (4.33)$$

$$\epsilon^B = \frac{ge^{\frac{\mu}{T}}T^4}{2\pi^2}(3a^2K_2(a) + a^3K_1(a)) \quad (4.34)$$

$$\begin{aligned} \epsilon^1 = & \frac{ge^{\frac{\mu}{T}}T^4}{4\pi^2} [9a^3K_3(a) + 4a^4K_2(a) + a^5K_1(a) \\ & + 2b(3a^2K_2(a) + a^3K_1(a) - 3a^3K_3(a) + a^4K_2(a)) \\ & + b^2(3a^2K_2(a) + a^3K_1(a))] \end{aligned} \quad (4.35)$$

In Fig. 4.13 we show the ratio of the energy density to first order in $(q - 1)$ to the full energy density as given by the Tsallis distribution, $(\epsilon^B + (q - 1)\epsilon^1)/\epsilon$ for several values of q indicated in the figure as a function of the temperature T .

Again, as noted previously for the particle density, it can be seen that the expansion in $(q - 1)$ is excellent if $(q - 1) = 0.01$ but rapidly deviates from the full Tsallis distribution for larger values of q . Also here, for $(q - 1) \approx 0.1$ the deviations are of the order of 20 % as can be seen from Fig. 4.13.

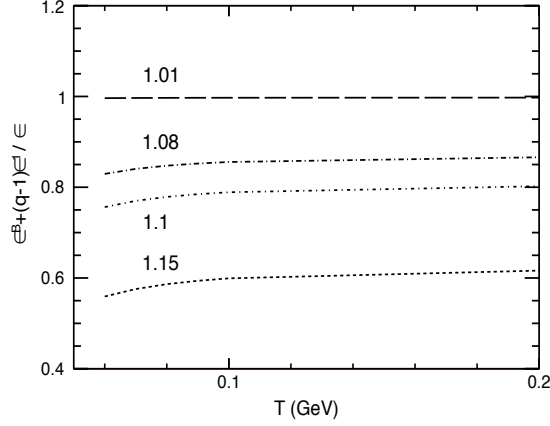


Figure 4.13: The ratio of the energy density calculated to first order in $(q - 1)$ normalized to the energy density of a Tsallis gas as a function of the temperature for different values of the parameter q . The mass is taken as being the pion mass. The values of the parameter q are 1.01 for the dashed line, 1.08 for the dot-(long)dashed line, 1.1 for the dot-(short)dashed line and 1.15 for the dotted line [95].

For comparison we show in Fig. 4.14, the first order expansion compared to the Boltzmann expression, $(\epsilon^B + (q - 1)n^1)/\epsilon^B$, as a function of the temperature T for several values of the parameter q . As in the previous case the deviations are most pronounced for small values of the temperature. Finally, the pressure is given by

$$P^B + (q - 1)P^1 \quad (4.36)$$

$$P^B = \frac{ge^{\frac{\mu}{T}}T^4a^2K_2(a)}{2\pi^2} \quad (4.37)$$

$$P^1 = \frac{ge^{\frac{\mu}{T}}T^4}{4\pi^2} [a^4K_2(a) + 3a^3K_3(a) - 2a^3bK_3(a) + a^2b^2K_2(a) + 2a^2bK_2(a)] \quad (4.38)$$

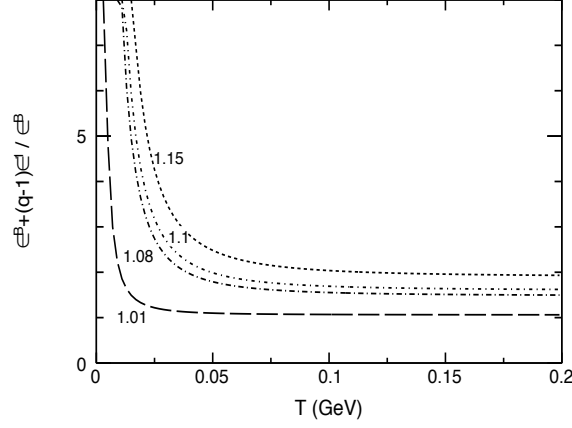


Figure 4.14: The ratio of the energy density calculated to first order in $(q - 1)$ normalized to the energy density of a Boltzmann gas as a function of the temperature for different values of the parameter q [95].

In Fig. 4.15 we show the ratio of the pressure to first order in $(q - 1)$ to the full pressure as given by the Tsallis distribution, $(P^B + (q - 1)P^1)/P$ for several values of q indicated in the figure as a function of the temperature T . Again, as noted previously, it can be seen that the expansion in $(q - 1)$ is excellent if $(q - 1) = 0.01$ but rapidly deviates from the full Tsallis distribution for larger values of q . Also here, for $(q - 1) \approx 0.1$ the deviations are of the order of 20 % as can be seen from Fig. 4.15. Again, we show in Fig. 4.16, the first order

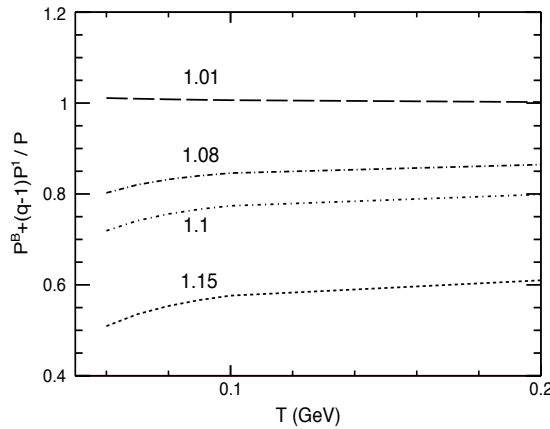


Figure 4.15: The ratio of the pressure calculated to first order in $(q - 1)$ normalized to the pressure of a Tsallis gas as a function of the temperature for different values of the parameter q [95].

expansion compared to the Boltzmann expression, $(P^B + (q - 1)P^1)/P^B$, as a function of the temperature T for several values of the parameter q . In this case the deviations are most pronounced for small values of the temperature.

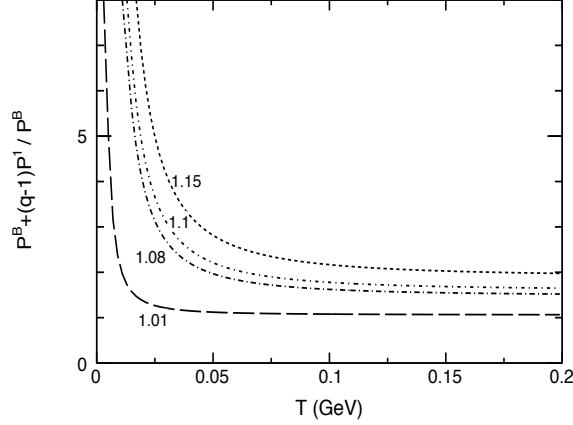


Figure 4.16: The ratio of the pressure calculated to first order in $(q - 1)$ normalized to the pressure of a Boltzmann gas as a function of the temperature for different values of the parameter q [95].

The entropy density can be obtained from the above expressions by using the thermodynamic relation

$$\epsilon + P = Ts + \mu n \quad (4.39)$$

and is not being shown here.

Before closing the section, we comment about the validity of the present expansion up to the second order in $(q - 1)$. As we have already seen, for truncation of the expansion up to the first order with pure Tsallis distribution, one needs to satisfy two conditions, i.e.

$$|1 - q| \frac{E}{T} < 1 \quad (4.40)$$

and

$$|1 - q| \left(\frac{E}{T} \right)^2 < 2 \quad (4.41)$$

With the thermodynamically consistent Tsallis distribution, the second condition turns out to be $|q(1 - q)|(E/T)^2 < 2$. For expansion up to second order in $(q - 1)$ with modified Tsallis distribution, the condition becomes $q^2|1 - q|(E/T)^3 < 3$.

Given the two values (the highest and the lowest) of $(q - 1)$ used in the present analysis, we want to put an upper bound in E/T until which the expansion will be reliable. If $(q - 1) = 0.15$, Eq.4.40 gives $E/T < 6.7$ and the modified condition for expansion up to first order gives $(E/T)^2 < 11.59$. That means for $T = 0.1$ GeV, the expansion will be reliable up to $E \approx 0.3$ GeV when $(q - 1) = 0.15$.

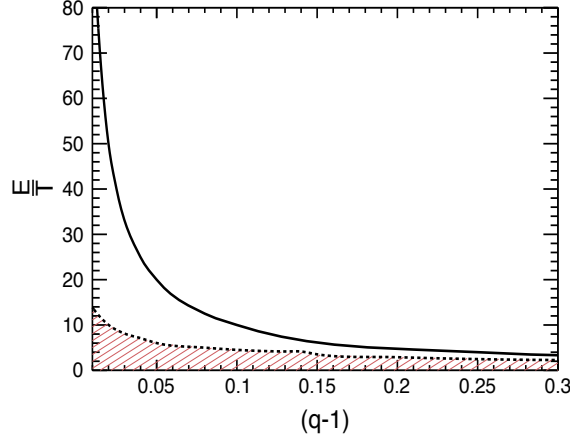


Figure 4.17: (Colour online) The region of validity for the expansion in $(q - 1)$ up to second order. The area under the solid line denotes the region where only the condition $|1 - q|E/T < 1$ is satisfied. The common overlapped area (i.e. the area under the dotted line) depicts the region within which all the conditions for expansion up to second order are satisfied [95].

If $(q - 1) = 0.01$, Eq.4.40 gives $E/T < 100$ and the modified condition for expansion up to first order gives $(E/T)^2 < 198$. That means for $T = 0.1$ GeV, the expansion will be reliable up to $E \approx 1.4$ GeV when $(q - 1) = 0.01$. The permissible values of $(q - 1)$ and E/T for reliable expansion up to the second order in $(q - 1)$ are shown to reside inside the smaller area filled with slanted lines. This is shown in Fig. 4.17. Hence, we conclude from the above discussion that the smaller the q value the more reliable the expansion becomes.

4.4.3 Inclusion of Flow to First Order in $(q - 1)$

In order to see how the inclusion of flow could improve the description of the transverse momentum distributions obtained in Pb-Pb collisions, we have included a

constant flow velocity, β . Assuming space-like freeze-out surface, the invariant yield is given by (see Appendix for the derivation).

$$\begin{aligned}
\frac{1}{p_T} \frac{dN}{dp_T dy} &= \frac{gV}{(2\pi)^2} \\
&\left\{ 2T[rI_0(s)K_1(r) - sI_1(s)K_0(r)] \right. \\
&\quad - (q-1)Tr^2I_0(s)[K_0(r) + K_2(r)] \\
&\quad + 4(q-1)TrsI_1(s)K_1(r) \\
&\quad - (q-1)Ts^2K_0(r)[I_0(s) + I_2(s)] \\
&\quad + \frac{(q-1)}{4}Tr^3I_0(s)[K_3(r) + 3K_1(r)] \\
&\quad - \frac{3(q-1)}{2}Tr^2s[K_2(r) + K_0(r)]I_1(s) \\
&\quad + \frac{3(q-1)}{2}Ts^2r[I_0(s) + I_2(s)]K_1(r) \\
&\quad \left. - \frac{(q-1)}{4}Ts^3[I_3(s) + 3I_1(s)]K_0(r) \right\} \quad (4.42)
\end{aligned}$$

where

$$r \equiv \frac{\gamma m_T}{T} \quad (4.43)$$

$$s \equiv \frac{\gamma v p_T}{T} \quad (4.44)$$

$I_n(s)$ and $K_n(r)$ are the modified Bessel functions of the first and second kind. Now, in this formula, the freeze-out surface has been considered to be space-like and so the integration over the freeze-out surface turns out to be trivial. For a more detailed treatment of the freeze-out surface in this context, the readers are referred to Ref. [99].

The comparison between model and experiment is quite good with notable deviations at small values of the transverse momentum p_T and again above values of 2.5 GeV (see Fig. 4.18).

Understanding the equation of state (EoS) of the system produced in high energy collisions is challenging and can be understood by looking into the speed of sound in that medium. This also helps in looking for any critical behaviour the system could undergo during its space-time evolution. The study of the speed of sound in hadronic matter hence becomes a subject of paramount importance in heavy-ion collisions.

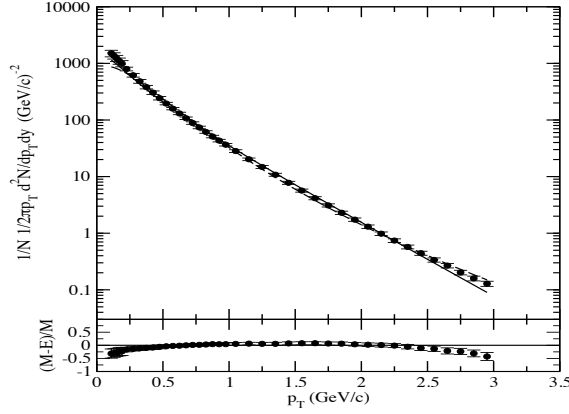


Figure 4.18: Fits to the normalized differential π^- yields as measured by the ALICE collaboration in $(0 - 5)\%$ Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [96]. The fit with the Tsallis distribution including flow keeping terms to first order in $(q - 1)$ (dashed line). The flow velocity is fixed at $\beta = 0.609$, with $T = 146$ MeV, $q = 1.030$ and the radius of the volume is $R = 29.8$ fm. The solid line is the Tsallis distribution without flow as given in Fig. 4.10. The lower part of the figure shows the difference between model (M), i.e. Tsallis with flow up to first order in $(q - 1)$, and experiment (E) normalized to the model (M) values [95].

4.5 Speed of Sound in Hadronic matter using Non-extensive Tsallis Statistics

The speed of sound (c_s) is studied to understand the hydrodynamical evolution of the matter created in heavy-ion collisions. The quark-gluon plasma (QGP) formed in heavy-ion collisions evolves from an initial QGP to the hadronic phase via a possible mixed phase. Due to the system expansion in a first order phase transition scenario, the speed of sound reduces to zero as the specific heat diverges. This also helps in looking for any critical behaviour the system could undergo during its space-time evolution. The study of the speed of sound in hadronic matter hence becomes a subject of paramount importance in heavy-ion collisions. The temperature dependence of the speed of sound in a medium is well established but the effect of temperature fluctuations is less explored, particularly in the case of high-energy collisions, where the non-extensive parameter q lies in the

range: $1 < q < 1.2$ [93]. As $q = 1$ corresponds to an equilibrated system described by Boltzmann-Gibbs statistics, the systems with small values of $(q - 1)$ are assumed to undergo an infinitesimal perturbation and tend to approach equilibrium. In non-extensive statistics, the Tsallis parameter (q) is related to the temperature fluctuations [100] and therefore we estimate the speed of sound using Tsallis statistics in hadronic medium, and study it as a function of q , for systems undergoing small perturbations. Here, we study the speed of sound (c_s) for a physical hadron resonance gas using the Tsallis non-extensive statistics to explore the underlying physical mechanisms.

4.5.1 Thermodynamics of a Hadron Resonance Gas

As we make a direct comparison of the thermodynamic quantities with those of an ideal pion gas, obeying Boltzmann statistics, for completeness we derive the necessary tools following Ref. [101]. The grand canonical partition function for an ideal pion gas having three degrees of freedom is given by

$$Z(T, V) = \sum_N \frac{1}{N!} \left[\frac{3V}{(2\pi)^3} \int_0^\infty d^3p \exp \frac{-\sqrt{p^2 + m_0^2}}{T} \right]^N, \quad (4.45)$$

where V is the system volume at temperature T and m_0 is the mass of a pion. The logarithm of the partition function gives

$$\ln Z(T, V) = \frac{3VTm_0^2}{2\pi^2} K_2(m_0/T), \quad (4.46)$$

One obtains the basic thermodynamic quantities like, pressure (P_0), energy density (ϵ_0) and entropy density (s_0) using the above partition function as

$$P_0(T) = T \left(\frac{\partial \ln Z}{\partial V} \right)_T = \frac{3m_0^2 T^2}{2\pi^2} K_2(m_0/T), \quad (4.47)$$

$$\begin{aligned} \epsilon_0(T) &= T^2 \left(\frac{\partial \ln Z}{\partial T} \right) \\ &= \frac{3m_0^2 T^2}{2\pi^2} [3K_2(m_0/T) + (m_0/T) K_1(m_0/T)], \end{aligned} \quad (4.48)$$

and

$$\begin{aligned} s_0(T) &= \frac{\epsilon_0(T) + P_0(T)}{T} \\ &= \frac{3m_0^2 T}{2\pi^2} [4K_2(m_0/T) + (m_0/T) K_1(m_0/T)], \end{aligned} \quad (4.49)$$

where K_1 and K_2 are the well-known modified Bessel functions of second kind.

For an ideal gas with zero chemical potential, the temperature dependent speed of sound, $c_s(T)$, is given by

$$c_s^2(T) = \left(\frac{\partial P}{\partial \epsilon} \right)_V = \frac{s(T)}{C_V(T)}, \quad (4.50)$$

where

$$s = \left(\frac{\partial P}{\partial T} \right)_V \quad (4.51)$$

is the entropy density and

$$C_V(T) = \left(\frac{\partial \epsilon}{\partial T} \right)_V \quad (4.52)$$

is the specific heat at constant volume. The specific heat at constant volume for an ideal pion gas using Eqs. 4.48 and 4.52 is given by

$$C_V^0(T) = 3s_0(T) + \frac{3m_0^4}{2\pi^2 T} K_2(m_0/T) \quad (4.53)$$

Using Eqn. 4.50 we get the speed of sound, c_s as

$$\frac{1}{c_s^2} - 3 = \frac{3m_0^4 K_2(m_0/T)}{2\pi^2 T s_0} \quad (4.54)$$

$$= \frac{m_0^2 K_2(m_0/T)}{4T^2 K_2(m_0/T) + m_0 T K_1(m_0/T)}, \quad (4.55)$$

with m_0 is the pion mass, the above expression gives the speed of sound for an ideal pion gas.

4.6 Speed of Sound in a Physical Hadron Resonance Gas

A hadron resonance gas consists of mesons and baryons obeying Bose-Einstein (BE) and Fermi-Dirac (FD) statistics, respectively. The Tsallis form of the Fermi-

Dirac and Bose-Einstein distributions as proposed in Refs. [102–107] is

$$f_T(E) \equiv \frac{1}{\exp_q\left(\frac{E-\mu}{T}\right) \pm 1}. \quad (4.56)$$

where the function $\exp_q(x)$ is defined as

$$\exp_q(x) \equiv \begin{cases} [1 + (q-1)x]^{1/(q-1)} & \text{if } x > 0 \\ [1 + (1-q)x]^{1/(1-q)} & \text{if } x \leq 0 \end{cases} \quad (4.57)$$

and, in the limit where $q \rightarrow 1$ it reduces to the standard exponential; $\lim_{q \rightarrow 1} \exp_q(x) \rightarrow \exp(x)$. In the present context we have taken $\mu = 0$, therefore $x \equiv E/T$ is always positive. In Eqn. 4.56, the negative sign in the denominator stands for BE and the positive stands for FD distribution.

To see the effect of quantum statistics on the speed of sound for a pion gas, we have explicitly compared the Tsallis-BE and Tsallis-Boltzmann distributions, which are shown in Fig. 4.19 as a function of temperature scaled with the Hagedorn limiting temperature, T_H . It is observed that for $q = 1.15$, there is no distinction between both the distributions, *i.e.* between the Tsallis-Boltzmann and the Tsallis-BE gas, and hence one considers Tsallis-Boltzmann distributions as a good approximation for describing the thermodynamics of a pion gas, as is done in the previous section. However, as expected, there is a clear distinction between equilibrium and Tsallis non-extensive statistics, for which the degree of deviation is higher towards lower temperatures.

For a hadron resonance gas, the energy density, pressure and specific heat at constant volume in non-extensive statistics are given by

$$\begin{aligned} \epsilon &= \frac{g}{2\pi^2} \int_0^\infty dp \, p^2 \sqrt{p^2 + m_0^2} \\ &\times \left[\left[1 + \frac{(q-1)\sqrt{p^2 + m_0^2}}{T} \right]^{\frac{1}{q-1}} \mp 1 \right]^{-q} \end{aligned} \quad (4.58)$$

$$\begin{aligned} P &= \frac{g}{2\pi^2} \int_0^\infty dp \, p^4 \frac{1}{3\sqrt{p^2 + m_0^2}} \\ &\times \left[\left[1 + \frac{(q-1)\sqrt{p^2 + m_0^2}}{T} \right]^{\frac{1}{q-1}} \mp 1 \right]^{-q} \end{aligned} \quad (4.59)$$

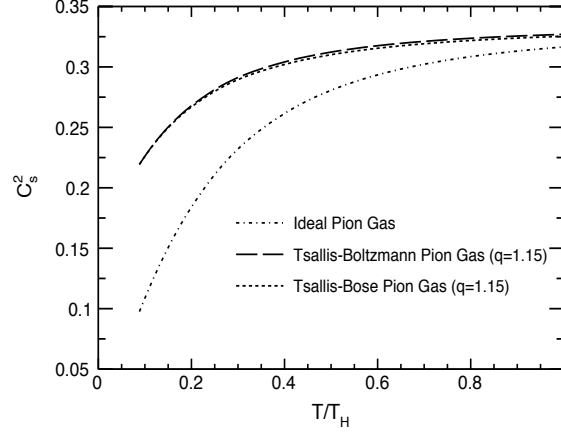


Figure 4.19: Speed of sound as a function of scaled temperature for an ideal pion gas (dash-dotted line: Boltzmann distribution) and for a pion gas with non-extensive statistics using Boltzmann (long-dashed line) and Bose-Einstein statistics (dotted line), for $q = 1.15$ [108].

and,

$$C_V = \frac{qg}{2\pi^2 T^2} \int_0^\infty dp p^2 (p^2 + m_0^2) \left[1 + \frac{(q-1)\sqrt{p^2 + m_0^2}}{T} \right]^{\frac{2-q}{q-1}} \times \frac{1}{\left[\left[1 + \frac{(q-1)\sqrt{p^2 + m_0^2}}{T} \right]^{\frac{1}{q-1}} \mp 1 \right]^{q+1}} \quad (4.60)$$

Hence the speed of sound can be calculated from these relations using Eqn. 4.50.

Figure 4.20 shows the speed of sound for a hadron resonance gas and the effect of taking different mass cut-offs of hadrons for a very small value of q ($=1.005$). The mass cut-off M is introduced as the highest mass of the resonances contributing to the hadron resonance gas. Ref. [109] has also considered the case of contributions going beyond such a cut-off. It is evident from Fig. 4.20 that in the limit of $q \sim 1$, a hadron resonance gas in Tsallis framework assuming quantum statistics gives identical results as shown in Ref. [101] for a similar system.

An increase in the value of q above one means a deviation from equilibrium statistics. To study the effect of higher values of q on the speed of sound, we have taken $q = 1.05$ and 1.15 with different mass cut-offs by inclusion of heavy

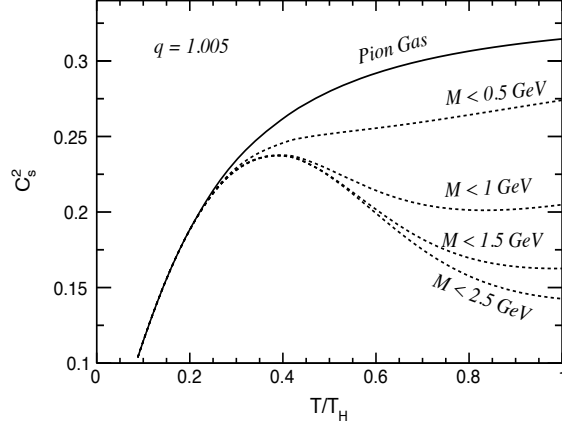


Figure 4.20: Speed of sound for $q=1.005$ for a hadron resonance gas with different cut-off on mass [108].

resonances.

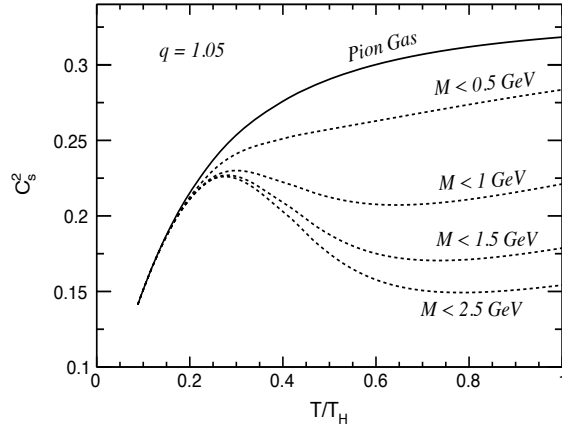


Figure 4.21: Speed of sound for $q=1.05$ for a hadron resonance gas with different cut-off on mass [108].

Figures 4.20 and 4.21 show that, for small values of q and large mass cut-off, the speed of sound at first increases as a function of temperature, reaches a maximum value and then decreases, which is indicative of a phase transition. In a first order phase transition it is expected that the speed of sound drops to zero at the phase transition point. This behavior is no longer present for larger values of q as can be seen from Figure 4.22 where the speed of sound increases monotonically with temperature. This basically is the behavior for any value of the mass cut-off. A comparison of Figures 4.20 and 4.21 with Figure 4.22 opens

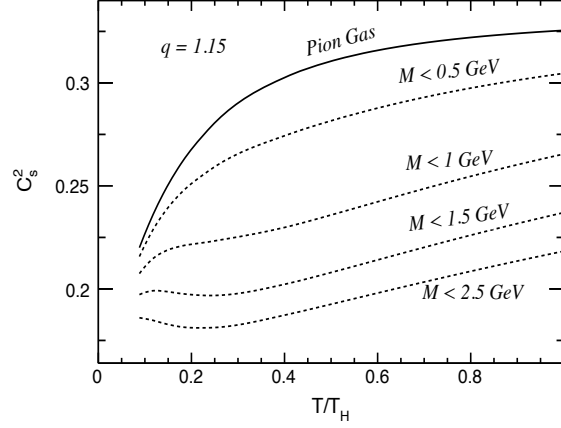


Figure 4.22: Speed of sound for $q=1.15$ for a hadron resonance gas with different cut-off on mass [108].

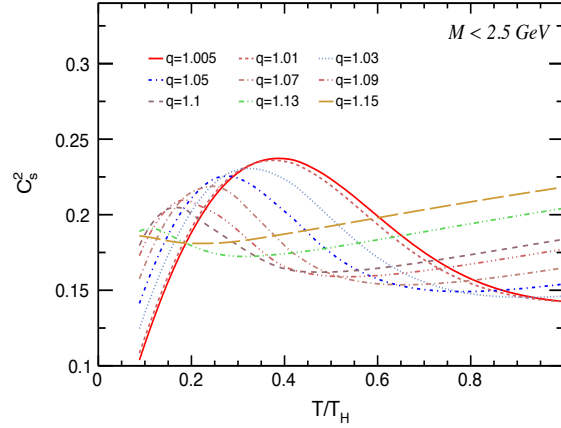


Figure 4.23: (Color online) Speed of sound for different values of the non-extensivity parameter, q , for a mass cut-off of $M < 2.5$ GeV [108].

the possibility that there exists a special value of q where a phase transition is no longer present. To explore this, we have studied the square of the speed of sound as a function of the scaled temperature for different values of q , taking a mass cut-off of 2.5 GeV. This is shown in Fig. 4.23. It could be inferred from the figure that with the progressive increase of the q -values, when the system goes away from equilibrium, the speed of sound slowly decreases to a minimum value up to $q = 1.13$. This behaviour, however, vanishes for higher values of q . This indicates the softening of the equation of state and a possible phase transition. Interestingly, this critical value of q is close to the value one obtains in the analysis

of p_T spectra in $p + p$ collisions at high energies.

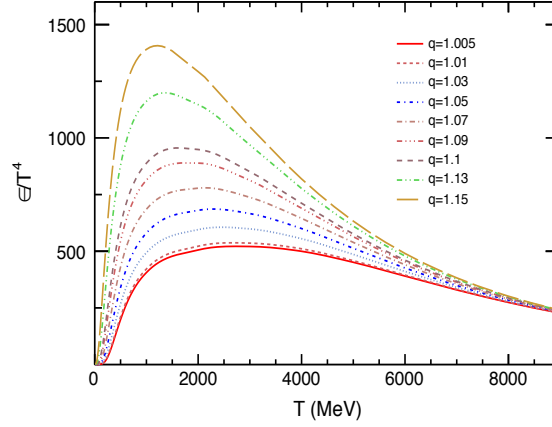


Figure 4.24: (Color online) Energy density scaled to T^4 (representing the degrees of freedom of the system) for different values of the non-extensivity parameter, q , for a mass cut-off of $M < 2.5$ GeV [108].

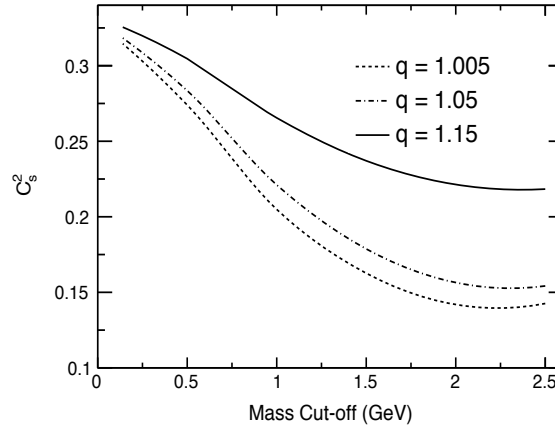


Figure 4.25: Speed of sound as a function of upper mass cut-off for a hadron resonance gas at $T = 170$ MeV, with different values of the parameter q [108].

In order to further explore the nature of the phase transition, we have studied the energy density (ϵ) scaled with T^4 , which is a measure of the degrees of freedom of the system, as a function of the temperature, taking the same mass cut-off for the physical resonance gas. This is shown in Fig. 4.24. With the increase of the q -values, ϵ/T^4 increases faster, showing a consistent sharp increase in its value for some temperatures, which may indicate a possible transition to

higher degrees of freedom. Then after reaching a maximum value for all q -values, it starts decreasing for higher temperature and goes to a saturation region. This behavior can be understood as follows. As the temperature increases more and more heavy resonances start contributing to the energy density and hence one notices a strong increase in the energy density. At infinite temperature, the system starts behaving like a gas of massless particles and ϵ/T^4 ultimately goes to a constant value. The rapid increase with the parameter q occurs because the distribution deviates more and more from a Boltzmann distribution and the tails at large momentum contribute more and more as q increases. For higher q -values the critical behaviour of c_s^2 diminishes by including heavy resonances. The speed of sound as a function of upper mass cut-off for a hadron resonance gas at temperature, $T = 170$ MeV, with different q -values is shown in Fig. 4.25. We choose a limiting temperature, $T_H \sim T_c = 170$ MeV motivated by the lattice QCD calculations [110] and also the Hagedorn limiting temperature [111–113], where T_H is the Hagedorn limiting temperature and T_c is the critical temperature for a deconfinement transition. This shows a monotonic decrease of speed of sound with the inclusion of higher resonances to the system and a systematic trend in the increase of the speed of sound is seen with the increase in q -value of the system.

Again experiments like ultra-relativistic heavy-ion colliders such as Relativistic Heavy Ion Collider (RHIC) and Large Hadron Collider (LHC) aim to produce matter at extreme conditions of temperature and/or energy density, where a phase transition is expected to take place from colourless hadronic matter to a coloured phase of quarks and gluons known as Quark-Gluon Plasma (QGP). Lattice QCD (lQCD) predicts a smooth crossover transition between hadron gas (HG) phase and QGP phase at zero baryon chemical potential [114, 115]. However, various QCD inspired phenomenological models predict a first order phase transition from HG to QGP phase [116–119]. These results advocate for a possible presence of a critical point (CP), where the first order phase transition ends. Currently, theoretical and experimental studies are dedicated towards understanding the nature of the QCD phase transition and location of CP. Thermodynamic observables such as specific heat and isothermal compressibility are very useful

in quantifying the nature of the phase transition and to obtain the equation of state (EOS) of the matter. Specific heat measures the change in energy density with respect to the change in temperature and is also related to the temperature fluctuation in the system [22–24]. The isothermal compressibility (κ_T) defines the change in volume with the change in pressure at constant temperature. It can be expressed as a second order derivative of Gibbs free energy with respect to pressure and is expected to diverge at CP. Thus, κ_T and C_V are very helpful in unveiling the nature of the phase transition and helpful in the search for CP. In the following work, we will focus on the isothermal compressibility (κ_T).

4.7 Isothermal Compressibility (κ_T)

Here, we present the formulations used to calculate the isothermal compressibility (κ_T) for a hot hadronic matter. First, we discuss hadron resonance gas (HRG) model used for calculation of κ_T . We then present the formulation of Excluded-Volume Hadron Resonance Gas model based on van der Waals type of interactions. Finally, we derive isothermal compressibility after including the Hagedorn mass spectrum in the HRG model.

4.7.1 Hadron Resonance Gas (HRG) Model

In this work, we consider a non-interacting physical hadron resonance gas. The grand canonical partition function for an ideal system of i^{th} particle species contained in a volume V in Boltzmann's approximation with mass m_i and chemical potential μ_i is given as [120],

$$\ln Z_i = \frac{g_i V}{2\pi^2} \int p^2 dp \exp [-(\varepsilon_i - \mu_i)/T], \quad (4.61)$$

where g_i is the degeneracy factor and $\varepsilon_i = \sqrt{p^2 + m_i^2}$ is the energy of i^{th} particle. The number density of hadrons calculated by using the basic thermodynamical relationship given as follows,

$$n_i = \frac{g_i}{2\pi^2} \int p^2 dp \exp [-(\varepsilon_i - \mu)/T]. \quad (4.62)$$

The isothermal compressibility is a measure of change of volume with the change in pressure at a constant temperature and average number of particles, which is given as [121],

$$\kappa_T \Big|_{T, \langle N_i \rangle} = -\frac{1}{V} \left(\frac{\partial V}{\partial P} \right) \Big|_{T, \langle N_i \rangle} \quad (4.63)$$

In order to calculate κ_T , we proceed as follows. The change in pressure can be written as,

$$dP = \left(\frac{\partial P}{\partial T} \right) dT + \sum_i \left(\frac{\partial P}{\partial \mu_i} \right) d\mu_i, \quad (4.64)$$

which gives,

$$\frac{\partial P}{\partial V} \Big|_{T, \langle N_i \rangle} = \sum_i \left(\frac{\partial P}{\partial \mu_i} \right) \left(\frac{\partial \mu_i}{\partial V} \right) \Big|_{T, \langle N_i \rangle}. \quad (4.65)$$

The first factor is the formula for N_i and the second factor is obtained from the condition of constancy of N_i as follows,

$$dN_i = \left(\frac{\partial N_i}{\partial T} \right) dT + \left(\frac{\partial N_i}{\partial V} \right) dV + \left(\frac{\partial N_i}{\partial \mu_i} \right) d\mu_i. \quad (4.66)$$

If N_i and T are kept constant, then Eq. 4.66 reduces to

$$\left(\frac{\partial \mu_i}{\partial V} \right) \Big|_{T, \langle N_i \rangle} = - \frac{\left(\frac{\partial N_i}{\partial V} \right)}{\left(\frac{\partial N_i}{\partial \mu_i} \right)}. \quad (4.67)$$

In the HRG model, $\frac{\partial N}{\partial V} = \frac{\partial P}{\partial \mu}$. Thus, Eq. 4.65 becomes,

$$\left(\frac{\partial P}{\partial V} \right) \Big|_{T, \langle N_i \rangle} = - \sum_i \frac{\left(\frac{\partial P}{\partial \mu_i} \right)^2}{\left(\frac{\partial N_i}{\partial \mu_i} \right)}. \quad (4.68)$$

Using Eqs. 4.63 and 4.68, we get,

$$k_T \Big|_{T, \langle N_i \rangle} = \frac{1}{\sum_i \left[\left(\frac{\partial P}{\partial \mu_i} \right)^2 / \left(\frac{\partial N_i}{\partial \mu_i} \right) \right]}. \quad (4.69)$$

Here, $n_i = N_i/V = \frac{\partial P}{\partial \mu_i}$ and calculated by using equation 4.62.

Eq. 4.69 gives the expression for the estimation of isothermal compressibility (κ_T). Let's now consider an excluded-volume HRG for studying isothermal compressibility (κ_T) from lower centre of mass energies to the highest LHC energies. EV-HRG is expected to give different results at lower $\sqrt{s_{NN}}$, where the assigned hard-core size of a baryonic volume plays an important role.

4.7.2 Excluded-Volume Hadron Resonance Gas (EV-HRG) Model

Now, we describe an equation of state for the HG based on the excluded-volume correction [122], where a hard-core size is assigned to all the baryons while mesons are treated as point-like particles in the grand canonical partition function. An excluded-volume correction in physical hadron resonance gas indirectly takes care of van der Waals-like repulsive interactions that arise due to a hard core size of the particles. In this approach, an annihilation of baryons and anti-baryons is not allowed. The excluded number density is calculated as [122],

$$n_i^{ex} = \frac{\lambda_i}{V} \left(\frac{\partial \ln Z_i^{ex}}{\partial \lambda_i} \right)_{T,V}, \quad (4.70)$$

with

$$\ln Z_i^{ex} = V(1 - \sum_j n_j^{ex} V_j^0) I_i \lambda_i, \quad (4.71)$$

Here, V_j^0 is the eigen-volume assigned to each baryon of j^{th} species and hence $\sum_j n_j V_j^0$ becomes the total occupied volume where n_j represents the total number of baryons of j^{th} species.

Equation 4.70 leads to a transcendental equation:

$$n_i^{ex} = (1 - R) I_i \lambda_i - I_i \lambda_i^2 \frac{\partial R}{\partial \lambda_i} + \lambda_i^2 (1 - R) I_i', \quad (4.72)$$

where I_i' is the partial derivative of I_i with respect to λ_i and $R = \sum_i n_i^{ex} V_i^0$ is the fractional occupied volume. We can write R in an operator equation form as follows [123]:

$$R = R_1 + \hat{\Omega} R, \quad (4.73)$$

where $R_1 = \frac{R^0}{1+R^0}$ with $R^0 = \sum n_i^0 V_i^0 + \sum I_i' V_i^0 \lambda_i^2$; n_i^0 is the density of point-like baryons of i^{th} species and the operator $\hat{\Omega}$ has the form:

$$\hat{\Omega} = -\frac{1}{1+R^0} \sum_i n_i^0 V_i^0 \lambda_i \frac{\partial}{\partial \lambda_i}. \quad (4.74)$$

Using Neumann iteration method and retaining the series upto $\hat{\Omega}^2$ term, we get

$$R = R_1 + \hat{\Omega} R_1 + \hat{\Omega}^2 R_1. \quad (4.75)$$

Once we get R , we can calculate the excluded number density and isothermal compressibility by using equations 4.69 and 4.72, respectively.

At higher collision energies, one expects for the production of high mass resonances. Theoretically, Hagedorn mass spectrum gives an opportunity to study the effect of higher resonances on various observables to explore QCD phase diagram in focus with the search for a CP.

4.7.3 Hagedorn Mass Spectrum

In this subsection, we include exponentially increasing continuous mass spectrum of the Hagedorn form given by eq. 4.10 in the grand canonical partition function for a hadron gas, which is given as,

$$\begin{aligned} \ln Z = \frac{1}{2\pi^2} \int \exp\left(-(\varepsilon - \mu)/T\right) p^2 dp \\ \times \int_{m_0}^{\infty} \rho(m) dm, \end{aligned} \quad (4.76)$$

where, the spectral function, $\rho(m)$ can be written as,

$$\rho(m) = \frac{A}{(m^2 + m_0^2)^{5/4}} \exp(m/T_H). \quad (4.77)$$

Now, eq. 4.76 becomes,

$$\begin{aligned} \ln Z = \frac{1}{2\pi^2} \int \exp\left(-(\varepsilon - \mu)/T\right) p^2 dp \\ \times \int_{m_0}^{\infty} \frac{A}{(m^2 + m_0^2)^{5/4}} \exp(m/T_H) dm. \end{aligned} \quad (4.78)$$

The values of A and m_0 , those arise from fitting the lQCD data are $0.57 \text{ GeV}^{-3/2}$ and 0.425 GeV , respectively [124]. Here, the Hagedorn temperature, $T_H = 0.180$

GeV [124]. The formula for the number density and pressure of hadrons after including the Hagedorn mass spectrum is given as follows,

$$n = \frac{1}{2\pi^2} \int \exp\left(-(\varepsilon - \mu)/T\right) p^2 dp \times \int_{m_0}^{\infty} \frac{A}{(m^2 + m_0^2)^{5/4}} \exp(m/T_H) dm, \quad (4.79)$$

and,

$$P = \int \exp\left(-(\varepsilon - \mu)/T\right) \frac{p^4}{3\varepsilon} dp \times \int_{m_0}^{\infty} \frac{A}{(m^2 + m_0^2)^{5/4}} \exp(m/T_H) dm. \quad (4.80)$$

The isothermal compressibility of hadrons in the presence of Hagedorn mass spectrum is calculated by using eq. 4.69.

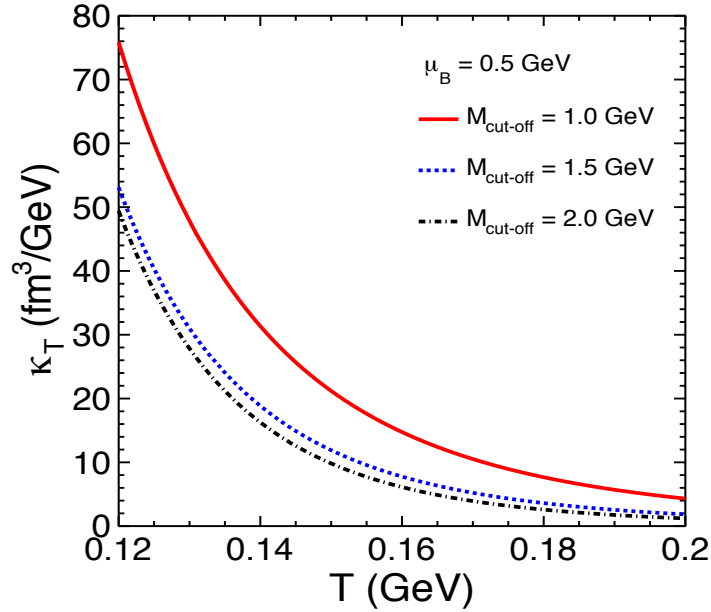


Figure 4.26: (Color online) The isothermal compressibility as a function of temperature for various mass cut-offs at $\mu_B = 0.5$ GeV [125].

In figure 4.26, we have shown the κ_T as a function of temperature for various mass cut-offs at $\mu_B = 0.5$ GeV calculated in HRG model. κ_T decreases with temperature due to increase of number of particles. We also observe that

isothermal compressibility has a lower value when higher mass resonances added into the system. This suggests that with more number of particles, it is difficult to compress the system and hence it becomes more incompressible. This could be because of particles gaining higher kinetic energy. Further, we have explicitly calculated the change in number density in HRG model with $\mu_B = 0.5$ GeV at $T = 0.170$ GeV for various mass cut-offs. We find 23% increase in number density when mass cut-off increases from 1.0 GeV to 1.5 GeV [125].

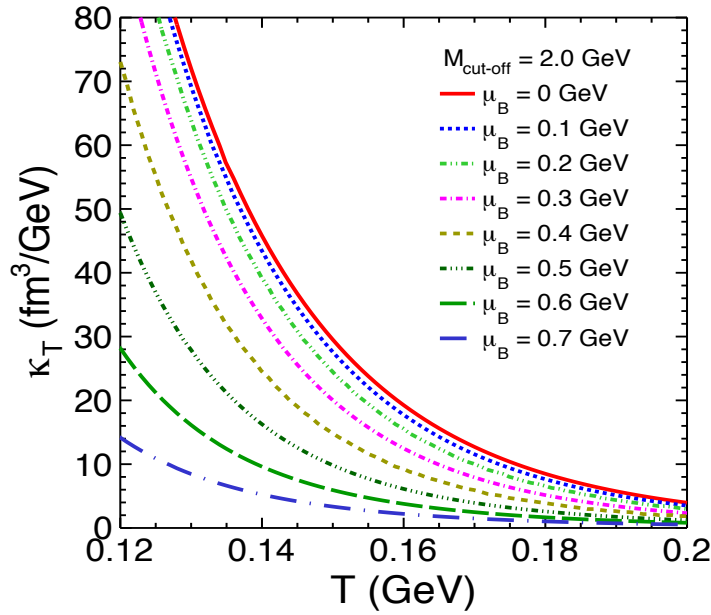


Figure 4.27: (Color online) κ_T as a function of temperature for various μ_B [125]. The lines are the results obtained in HRG model for different μ_B .

Figure 4.27 depicts the isothermal compressibility calculated in HRG model as a function of temperature for various μ_B with mass cut-off of 2 GeV. We notice that, κ_T decreases with the temperature and has lower values for higher baryon chemical potential, which goes inline with the concept that as the number of baryon states increase, the system becomes incompressible. The effect of baryon chemical potential on κ_T is more pronounced at a lower temperature.

In Fig. 4.28, a 3-dimensional (3D) plot of isothermal compressibility calculated in HRG model as a function of T and μ_B is shown. It is observed that κ_T decreases with T at a constant μ_B and its values decrease with the increasing μ_B

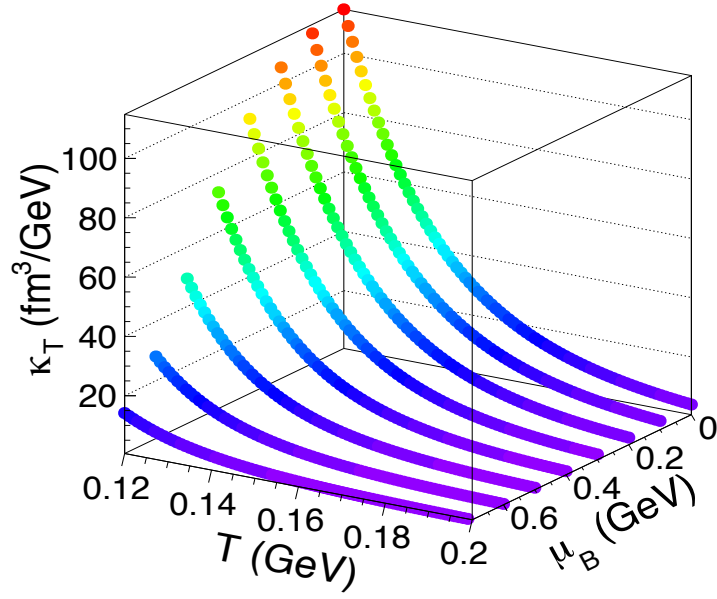


Figure 4.28: (Color online) Variation of κ_T as a function of T and μ_B [125].

for a given temperature of the system.

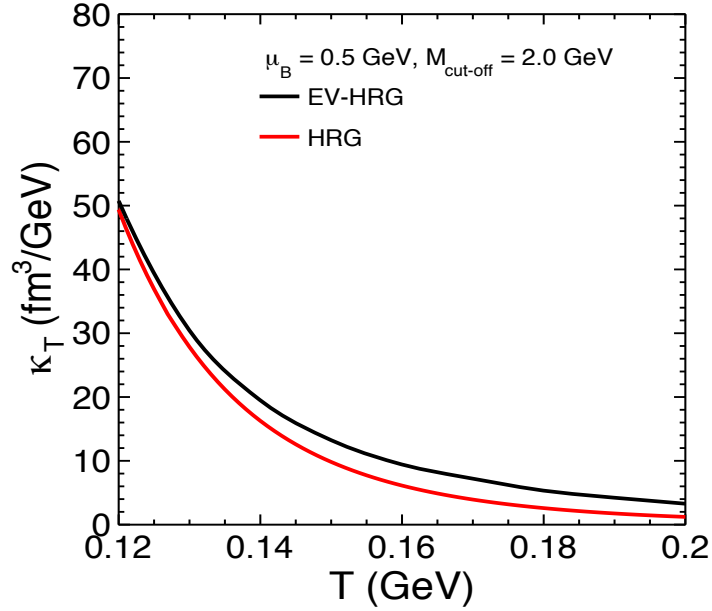


Figure 4.29: (Color online) κ_T as a function of temperature. Black line is the result obtained in EV-HRG and red line is from HRG calculations [125].

Now, we want to compare the results of HRG and EV-HRG models. For

this, we compare κ_T calculated in both the models for $\mu_B = 0.5$ GeV with mass cut-off = 2 GeV as shown in Fig. 4.29. We notice that when κ_T is varied as a function of T , the EV-HRG results lie above those obtained in HRG model and the difference is larger at a higher T . This finding tells that excluded-volume correction in ideal hadron gas model plays an important role as we go towards higher temperatures.

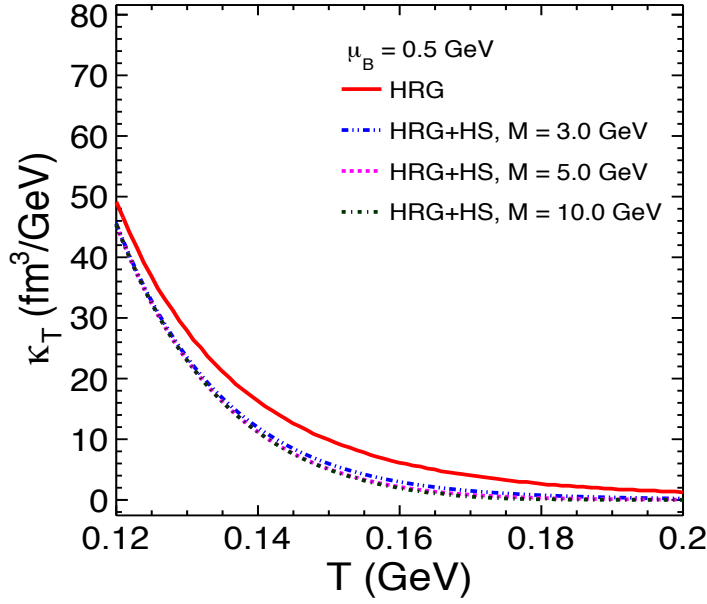


Figure 4.30: (Color online) κ_T as a function of temperature in HRG model and HRG with Hagedorn mass spectrum for various mass cut-offs at $\mu_B = 0.5$ GeV [125].

Figure 4.30 represents the results of κ_T estimated in the HRG model for $\mu_B = 0.5$ GeV. We have also shown the results after adding the Hagedorn states (HS) in the grand partition function of HRG model for various upper mass limits. Here, we take continuous mass spectrum in the calculation above hadron mass of 2 GeV. An upper mass limit in the mass integration is taken upto 10 GeV. The decrease in compressibility is found with the addition of higher HS into the system. The effect of Hagedorn states is prominent when the mass cut-off increases from 2 GeV to 3 GeV.

Figure 4.31 depicts the impact of baryochemical potential on the isother-

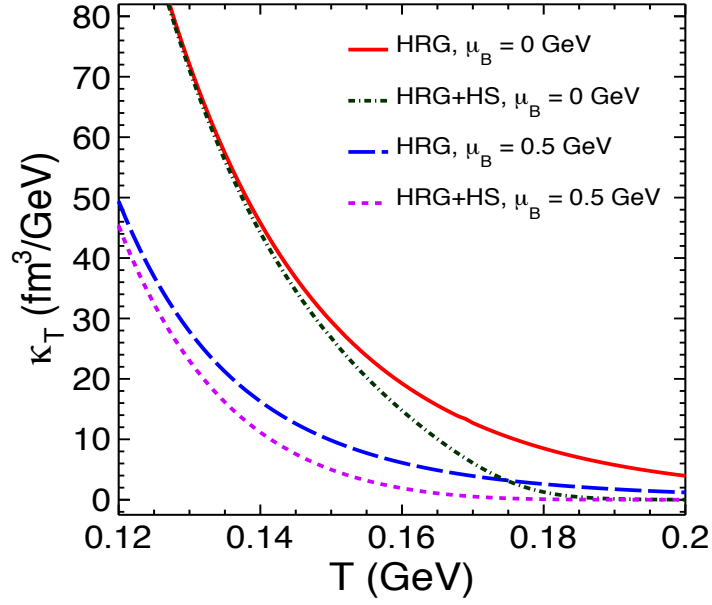


Figure 4.31: (Color online) κ_T as a function of temperature in HRG model and HRG with Hagedorn mass spectrum for various μ_B at mass cut-off of 10 GeV [125].

mal compressibility for hadronic matter calculated in both HRG and HRG+HS models, where an upper mass limit of 2 and 10 GeV are taken in the HRG and HRG+HS models, respectively. It is evident from the figure that κ_T has lower values particularly at a higher T in case of HRG+HS in comparison to that obtained in HRG. Again, we notice that it decreases with the increasing baryonic chemical potential for both the cases. Addition of Hagedorn states reduce the values of κ_T supports the earlier findings that inclusion of higher masses makes the system incompressible.

In order to study the collision energy dependence of isothermal compressibility, we need a parametrisation of temperature and baryon chemical potential in terms of centre-of-mass energy [6]:

$$\mu_B = \frac{a}{1 + b\sqrt{s_{NN}}}, \quad (4.81)$$

and,

$$T = c - d\mu_B^2 - e\mu_B^4, \quad (4.82)$$

where the parameters in EV-HRG are: $a = 1.482 \pm 0.004$ GeV, $b =$

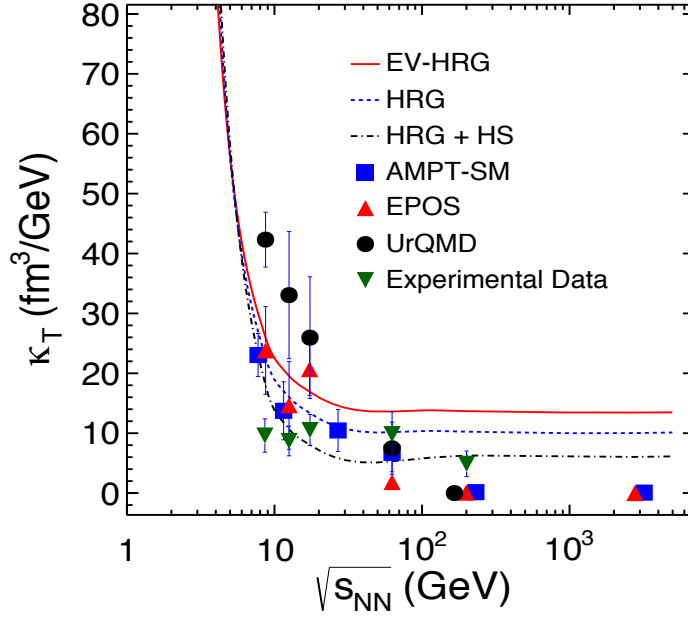


Figure 4.32: (Color online) The centre-of-mass energy dependence of κ_T obtained in excluded-volume, HRG and HRG with Hagedorn states [125]. The various transport models and experimental results shown by various symbols are taken from ref. [126].

$0.352 \pm 0.009 \text{ GeV}^{-1}$, $c = 0.163 \pm 0.002 \text{ GeV}$, $d = 0.17 \pm 0.02 \text{ GeV}^{-1}$ and $e = 0.015 \pm 0.010 \text{ GeV}^{-3}$. In case of HRG, the parameters are: $a = 1.22 \pm 0.04 \text{ GeV}$, $b = 0.242 \pm 0.017 \text{ GeV}^{-1}$, $c = 0.170 \pm 0.003 \text{ GeV}$, $d = 0.190 \pm 0.039 \text{ GeV}^{-1}$ and $e = 0.0108 \pm 0.0074 \text{ GeV}^{-3}$ [122]. Figure 4.32 represents κ_T as a function of centre-of-mass energy ($\sqrt{s_{NN}}$). The solid line is the result obtained in EV-HRG model while the dashed line is the result obtained in HRG model. The dash-dotted line is the calculation done in the framework of a HRG model with the Hagedorn mass spectrum. We perceive that isothermal compressibility initially decreases rapidly with $\sqrt{s_{NN}}$ and gets saturated around collision energy of 10-20 GeV. This change of behaviour of κ_T emphasises a very different nature of the system at lower collision energies in comparison to higher energies. The energy window at which such a change of behaviour of compressibility is observed could be the onset of phase transition from hadron gas to QGP. Significant differences are found between EV-HRG, HRG and HRG+HS results for $\sqrt{s_{NN}} > 10 \text{ GeV}$.

Further, we find that addition of Hagedorn states in the HRG model lowers down the values of isothermal compressibility and it is closer to the experimental data as compared to other model results. This necessitates the inclusion of such Hagedorn mass spectrum while studying thermodynamical observables. We have also shown the experimental results obtained in various collision energies and the results of transport models like AMPT with string melting scenario, UrQMD and EPOS [126]. AMPT with string melting scenario (AMPT-SM) and EPOS results are closer to the experimental data. However, UrQMD results deviate from experimental data at lower collision energies. This study helps in confining the search for critical point in the QCD phase diagram to a centre-of-mass energy of 10-20 GeV per nucleon. It is observed that the matter created in heavy-ion collisions at the lower collision energies like in future facilities e.g. CBM@FAIR, NICA, Dubna is more compressible, unlike that created at the RHIC and LHC energies. Similarly, Addition of Hagedorn mass spectrum in the HRG model reduces the isothermal compressibility and hence it approaches closer to experimental data at higher collision energies in comparison to the other theoretical calculations. This signifies the use of Hagedorn mass spectrum into the system while studying the isothermal compressibility.

4.8 Non-Extensive Statistics in Free-Electron Metals and Thermal Effective Mass

It would be exciting to explore possible applications of non-extensivity in condensed matter physics. Here, we use Tsallis thermodynamically consistent distribution in studying the thermal effective mass and specific heat of various metals. We have found an excellent agreement with experimental findings.

Study of properties of metals has great importance in condensed matter physics. Over two third of solids are metals and, they are good for both electrical and thermal conduction. Even to understand the non-metals we should understand the behaviour of metals. The Drude model, which applies the kinetic theory of gases to electrons in metals is successful to explain the basic transport properties in metals. It basically is based on the *independent* as well as *free-electron*

approximations ignoring the electromagnetic electron-electron and electron-ion interactions. Despite the oversimplification of the reality, it describes quite well the electrical and thermal conductivities of metals and the Hall effect [127]. However, Drude model fails to explain many other observations. According to Drude model the specific heat due to electrons in a metal is independent of temperature and can be represented as $C_V = 3nk_B/2$, where n is the electron density, k_B representing the Boltzmann constant. This shows metals with more number of free electrons should have much larger heat capacity. However, this is not experimentally observed and is the consequence of taking Maxwell-Boltzmann distribution to describe the energy distribution in a many-electrons system. Later, Sommerfeld applied the same principle to metals but with a modification to the electronic velocity (or momentum) distribution, which is taken to be a quantum FD distribution, derived from Boltzmann-Gibbs (BG) statistics. The Drude-Sommerfeld model remains valid and also allows us to understand many phenomena in metals despite the lack of accuracy and simplicity of its assumptions. But the key feature of this model is that electrons in a metal, as described by the FD distribution, are assumed to be in thermal equilibrium with the surroundings by collisions.

However, C_V , the specific heat due to the electrons, calculated using the above Drude-Sommerfeld model of free electrons is known to deviate from experimental data. In text books this deviation is usually explained in terms of a *thermal effective mass* of the electrons. In an actual situation, there are interactions and collisions (electron-electron, electron-ion etc) that can result in taking the system away from thermal equilibrium, at least locally. In such a case a non-extensive (NE) Tsallis statistics [128], instead of BG statistics, may be of use.

It must be mentioned here that BG statistics applies correctly to systems that satisfy broadly the following conditions. The interactions must be short ranged, Markovian type, boundary conditions must be smooth, and there must be no mesoscopic dissipation taking place. Violation of any or more of these condition would necessitate the application NE statistics as proposed by Tsallis.

In the present work [129], we have used non-extensive statistics to derive the electronic specific heat, C_V in some typical free-electron metals to provide an

explanation of the *thermal effective mass* of the electrons in terms of the non-extensive parameter q . The list of metals considered includes, Li, Na, Al, K, Ti, Fe, Co, Ni, Cu, Sr, Ag, Cs, Au, Hg and Pb, with 1, 2, 3 and 4 free electrons in the outer shell.

4.8.1 Finite Temperature Free Electron Gas and Specific Heat

At zero temperature, electrons would prefer the lowest energy states to minimize the energy. But following Pauli principle, they occupy energy states in pairs with opposite spins. If we have N electrons, at $T = 0$, they occupy the lowest $N/2$ states. The energy of the highest filled state at absolute zero, being the Fermi Energy, ϵ_F .

At zero temperature, the occupation probability $f(\epsilon)$, for a given quantum state with energy ϵ , is 1 if $\epsilon \leq \epsilon_F$ and 0 if $\epsilon \geq \epsilon_F$, leading to a step function

$$f(\epsilon) \equiv \begin{cases} 1 & \text{if } \epsilon \leq \epsilon_F \\ 0 & \text{if } \epsilon \geq \epsilon_F \end{cases} \quad (4.83)$$

At finite temperature interaction between electrons can be understood in terms of scattering leading to change in momentum or energy. Since all states with energy below ϵ_F are occupied scattering to low energy states are prohibited leading to a fraction of electrons with energy near ϵ_F getting excited to states above ϵ_F .

$$0 < f(\epsilon) < 1, \text{ at finite } T. \quad (4.84)$$

The total energy of a system within the free electron model is given by

$$U(T) = \int_0^\infty \epsilon f(\epsilon) D(\epsilon) d\epsilon \quad (4.85)$$

where

$$f(\epsilon) = \frac{1}{1 + \exp(\frac{\epsilon - \mu}{k_B T})} \quad (4.86)$$

is the FD distribution $f(\epsilon)$ for non-interacting electrons at finite temperature. Here μ and $D(\epsilon)$ are the chemical potential and density of states respectively. Since electrons near the Fermi energy take the extra energy, $D(\epsilon)$ is to be evaluated at $\epsilon = \epsilon_F$.

The heat capacity C_V , due to the electrons can be calculated using the following

$$C_V = \left. \frac{\partial U}{\partial T} \right|_V = - \int_0^\infty \frac{(\epsilon - \mu)^2}{T} \frac{df(\epsilon)}{d\epsilon} D(\epsilon) d\epsilon \quad (4.87)$$

where $D(\epsilon)$ is given by,

$$D(\epsilon) = \frac{dN}{d\epsilon} = \frac{3N}{2\epsilon_F}. \quad (4.88)$$

The specific heat can be expressed as

$$C_V = \frac{\pi^2}{2} N k_B \frac{k_B T}{\epsilon_F} \quad (4.89)$$

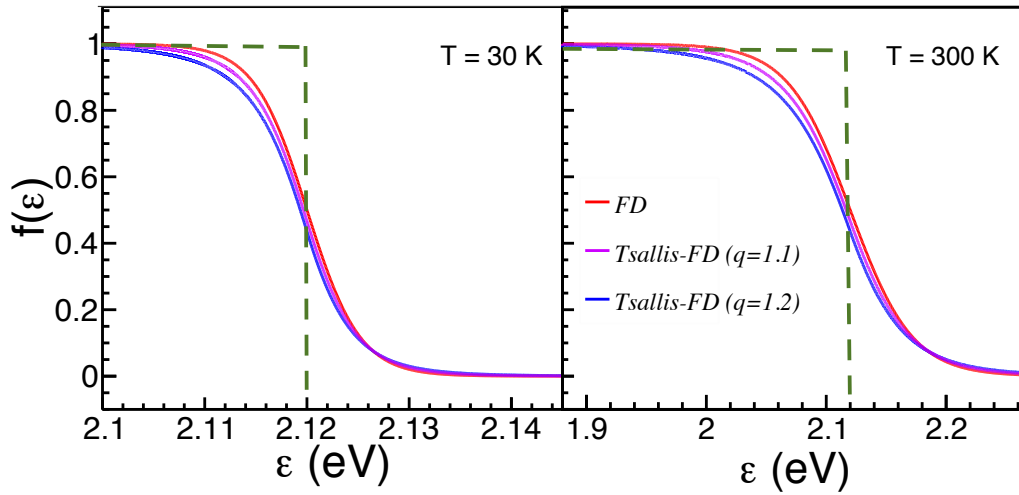


Figure 4.33: (Color online) Tsallis FD distribution of electrons in potassium with $\epsilon_F = 2.12 \text{ eV}$, for $q=1.1$ and 1.2 at two different temperatures, 30 K and 300 K. The dotted lines show the standard FD distribution for BG statistics at $T = 0$ for comparison [129]. For clarity all distributions are shown near the Fermi energy.

4.8.2 Tsallis Non-Extensive Statistics and Specific Heat

In the NE thermodynamics following Tsallis formalism, one uses a q -deformed quantum distribution using the q -exponential function (given a step later). The nonextensivity enters through the parameter q , usually close to unity. The relevant distribution in the present case is a modified FD distribution [75] as given by

$$f_T(\epsilon) = \frac{1}{1 + \exp_q\left(\frac{\epsilon - \mu}{k_B T}\right)}. \quad (4.90)$$

The Tsallis NE distribution function in the limit of the NE parameter, $q \rightarrow 1$ goes over to the standard equilibrium FD distribution function. The parameter q is a measure of the degree of non-extensivity based on the amount of correlations and possible temperature fluctuations that is not included in a pure free electron picture as described by the FD distribution. As has been mentioned earlier, in an actual case there are long range electron-ion and electron-electron correlations. In addition there could be local temperature fluctuations. However, it should be mentioned here that the q -statistics is phenomenological in nature and doesn't follow from the first principle. The use of NE statistics will be handy for such systems, where the Tsallis form of FD distribution as given in Eq. 4.90 could be used.

However, as shown elsewhere,[130, 131] just using the q -deformed distribution function is not enough to guarantee thermodynamic consistency. In other words, the formalism must comply with standard relationships between relevant thermodynamical variables such as entropy, temperature, energy etc. This requires two additional constraints $\sum n_i^q = N$ and $\sum n_i^q \epsilon_i = E$ where n_i and ϵ_i represent the mean occupation probability of the i^{th} quantum state with energy ϵ_i . E and N represent the total number of particles and the total energy respectively.

Using Tsallis NE statistics with the above mentioned constraints the total energy is given by,

$$U(T) = \int_0^\infty \epsilon f_T^q(\epsilon) D(\epsilon) d\epsilon \quad (4.91)$$

where the q^{th} power of the occupation number, $f_T(\epsilon)$, is used.

$$f_T^q(\epsilon) = \left(\frac{1}{1 + \exp_q\left(\frac{\epsilon - \mu}{T}\right)} \right)^q. \quad (4.92)$$

So, the heat capacity is

$$C_V = \frac{\partial U}{\partial T} \Big|_V = - \int_0^\infty \frac{(\epsilon - \mu)^2}{T} \frac{df_T^q(\epsilon)}{d\epsilon} D(\epsilon) d\epsilon \quad (4.93)$$

where $D(\epsilon)$ is the density of states at ϵ_F , given earlier in Eq. (9). C_V due to electrons can now be calculated from eqn. 4.93 using numerical techniques as used in the present case.

Table 4.2: Present list of metals, with atomic number, electron density, γ and the estimated q values. Numbers in the parentheses indicate references.

Elements	z	Electron-density ($10^{28}/\text{m}^3$)	ϵ_f (eV)	γ_0 [mJ/(mol K ²)]	γ_q [mJ/(mol K ²)]	γ_{Exp} [mJ/(mol K ²)]	m^*/m_0	q
Li	3	4.70	4.74	0.746	1.649	1.65 [132]	2.211	1.231
Na	11	2.65	3.24	1.091	1.381	1.38 [132]	1.265	1.082
Al	13	18.1	11.7	0.906	1.352	1.35 [132]	1.492	1.132
K	19	1.40	2.12	1.667	2.085	2.08 [132]	1.250	1.078
Ti	22	9.80	7.75	0.912	3.361	3.36 [132]	3.684	1.340
Fe	26	17.0	11.17	0.633	4.893	4.90 [132]	7.729	1.505
Co	27	15.63	10.58	0.668	4.410	4.40 [132]	6.599	1.467
Ni	28	18.2	11.74	0.602	7.032	7.04 [132]	11.676	1.617
Cu	29	8.47	7.0	0.505	0.690	0.69 [132]	1.366	1.106
Sr	38	3.55	3.95	1.790	3.645	3.64 [132]	2.036	1.212
Ag	47	5.86	5.49	0.644	0.648	0.64 [132]	1.006	1.0002
Cs	55	0.91	1.59	2.223	3.204	3.20 [133]	1.441	1.122
Au	79	5.90	5.53	0.639	0.692	0.69 [132]	1.083	1.028
Hg	80	8.65	7.13	0.992	1.865	1.86 [132]	1.880	1.193
Pb	82	13.2	9.37	1.509	2.996	2.99[132]	1.985	1.206

In Fig. 4.33, we show the FD distributions as a function of energy for both BG and Tsallis statistics, respectively for the case of potassium which is a nearly perfect free electron metal with an ϵ_F value of 2.12 eV. Results shown for the Tsallis statistics correspond to two different temperatures 30 K and 300 K. These results shown have been computed for two different q values 1.1 and 1.2. For clarity we show only results near the Fermi energy, ϵ_F where differences are very clearly seen. The corresponding result for the standard FD distribution coming from BG statistics, at $T = 0$, is also shown for comparison. One can clearly notice, at a given temperature, non-extensivity results in exciting more electrons to higher energy states, as compared to BG statistics based FD distribution. The net effect of this is an increased contribution to C_V . This effect becomes more and more important with increase in q .

4.8.3 Experimental Heat capacity in a real metal

In a real metal, at low temperature the total heat capacity can be more precisely written as

$$C_V = \gamma T + A T^3 \quad (4.94)$$

where γ and A are constants and related to the material. Here, the first linear term is the electronic contribution in which we are interested, the second cubic

term coming from lattice vibrations, based on the Debye model (phonons contribution). Therefore, at low temperatures, C_V/T when plotted against T^2 results in a straight line. In such a case the zero intercept corresponding to $T = 0$ yields the parameter γ . A compilation of this on a large number of systems is available [132, 133].

For a free electron gas, that obeys the FD distribution, we rename the coefficient γ as γ_0 . This is given by

$$\gamma_0 = \frac{\pi^2}{2} \frac{N k_B^2}{\epsilon_F} \quad (4.95)$$

where

$$\epsilon_F = \frac{\hbar^2}{2m} k_F^2 = \frac{\hbar^2}{2m} \left(3 \pi^2 \frac{N}{V} \right)^{2/3}. \quad (4.96)$$

So that

$$\gamma_0 = m \frac{\pi^2 N k_B^2}{\hbar^2 \left(3 \pi^2 \frac{N}{V} \right)^{2/3}} \quad (4.97)$$

One can now see, the parameter γ_0 is proportional to the free electron mass. Following the same prescription as valid for free electrons obeying the FD distribution, we can now determine γ for electrons following the Tsallis modified FD distribution as given in Eq. 4.90. We call it γ_q . This is possible provided one knows the value of the parameter q . In such a case the ratio γ_q/γ_0 can be seen as the ratio of the *thermal effective mass* of the electron, m^* , to free electron mass, m . In other words for a given system obeying the NE statistics, the parameter q can be determined from a knowledge of the m^*/m mass ratio. This is what has been followed in the present work.

In our analysis we first look at the case with potassium, a nearly perfect free electron system. In this case the total C_V including the lattice contribution (as given by Eq. (16)) has been obtained experimentally as a function of temperature, T [133]. In the low T region, C_V/T vs T^2 is shown to be perfectly linear. At $T = 0$, $C_V/T = \gamma$ corresponds to the intercept which has a value of 2.08 mJ/(mol.K²). Taking $\epsilon_F = 2.12$ eV as valid for potassium we estimate the parameter q to have a value of 1.078 which is greater than 1. In this estimation, varying the parameter q , we first determine C_V for a set of temperatures. Since there is no lattice contribution C_V vs T , in accordance with Eq. (15), is perfectly

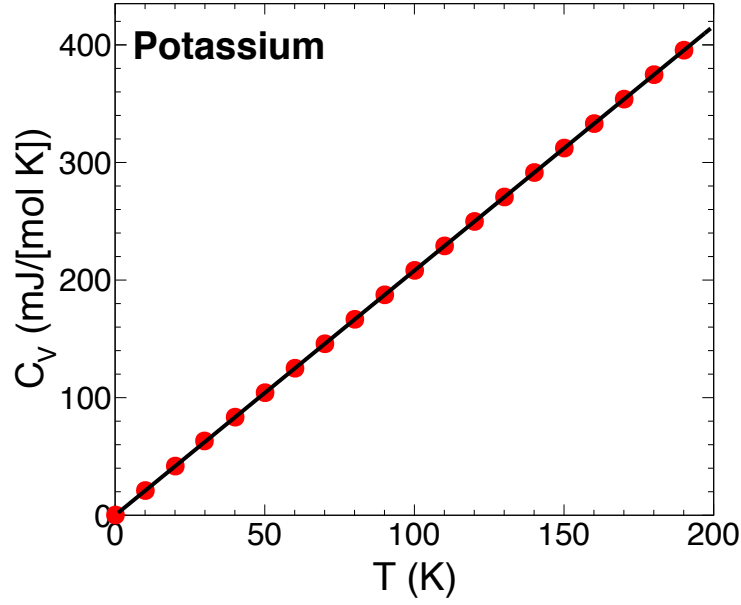


Figure 4.34: Electronic C_V versus T for potassium for $q=1.078$. The slope leads to a γ_q value of $2.085 \text{ mJ}/(\text{mol.K}^2)$ close to the experimental data [129].

linear with a zero intercept, the slope leading to γ_q . Our estimated C_V data for potassium, for q value of 1.078, at various temperatures is shown in Fig. 4.34. It has a slope (γ_q) value of 2.085, very close to the experimental data.

The same procedure is followed for the estimation of q and γ_q for other metals using the available data. As mentioned earlier, our full list includes Li, Na, Al, K, Ti, Fe, Co, Ni, Cu, Sr, Ag, Cs, Au, Hg and Pb, with 1, 2, 3 and 4 free electron systems as a complete set. The values, together with, other physical parameters such as electronic density, ϵ_F , γ_0 , γ_q and the ratio γ_q/γ_0 ($= m^*/m$) are presented in Table 4.2. In the same table m^*/m data and the experimental γ_{Expt} values available in the literature are also included for comparison. One can see, it is possible to adjust the NE parameter q to obtain γ_q values close to the experimental data.

From the results included in the table, one can further see, for all metals considered in this paper, with the exception of Ag, the non-extensive parameter, q deviates from unity, always being higher. All it means is that there are small local fluctuations in temperature coming from non-extensivity whose effect is taken care

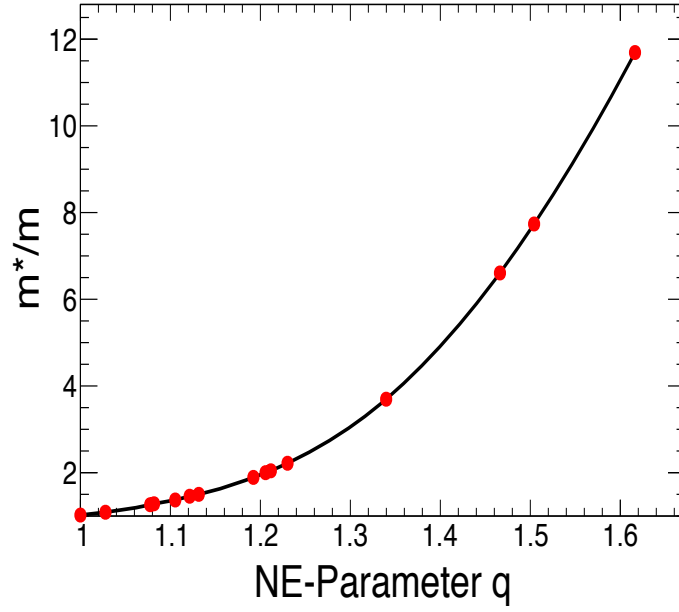


Figure 4.35: Ratio m^*/m for various metals with different q values. The data are well described by the the continuous line obtained using the reduced energy ϵ/ϵ_F as mentioned in the text [129].

of by the Tsallis statistics, leading to a higher value of γ ($= C_V/T$), resulting in a higher effective mass. It is important to point out that the parameter γ_q has been calculated in an earlier work [134], for a number of metals that include the present ones as well. Tsallis modified FD statistics was used in the calculations. where the total enegy was calculated using $f_T(\epsilon)$ rather than $f_T^q(\epsilon)$. As a result, the obtained γ values (shown in Table 1 of ref [134]) are not quite correct and are seen to be much less than the observed results. Even the values obtained for $q = 1$, corresponding to pure FD distribution, are in total disageement with free electron data, some of which are given in the present paper. This clearly emphasizes the need for thermodynamical consistency in the formalism which has been ensured in the present work.

At the end, in Fig.4.35 we show a plot of the ratio m^*/m as a function of the NE parameter q as obtained in various metals considered here. The results show a nice systematic dependence, all data corresponding to different values of ϵ_F indicating a smooth functional dependence. The parameter q distinguishes

one system from another as does ϵ_F . To check this further, for various values of q in the range 1-1.7, we repeated the calculations for each of the metals considered, changing the integration variable to ϵ/ϵ_F in the expression for C_V . In all the cases the scaled distribution functions fall off to zero near $\epsilon/\epsilon_F = 1$. This actually removes the dependence on ϵ_F bringing out the q dependence. The γ_q values obtained for each system were then divided by the corresponding γ_0 value leading to the m^*/m ratio as a function of q . Such a calculation, for each of the systems considered, results in the same smooth curve shown, passing exactly through all the points. This means, m^*/m for all free-electron-metals, irrespective of their Fermi energies, follow a universal trend in terms of the parameter q .

4.9 Summary

In this chapter, first we show the application of Tsallis non-extensive statistics to extract the information about the freeze-out surface of the identified particles in pp-collisions at $\sqrt{s} = 7$ TeV. Furthermore, we look for any possible hint of collective flow in high multiplicity pp collisions by analysing the transverse momentum spectra of the identified particles in pp-collisions at $\sqrt{s} = 7$ TeV using the BGBW model. We found Tsallis distribution provides a complete description of particle spectra upto very high- p_T where as BGBW explains the spectra up to 3 GeV/c. The variable T shows a systematic increase with multiplicity, the heaviest baryons showing the steepest increase. This is an indication of a mass hierarchy in particle freeze-out. The obtained parameters show variations with the event multiplicity. The notable variation of the non-extensive parameter, q which decreases towards the value 1 as the multiplicity increases and this effect is more significant for heavy mass particles. This shows the tendency of the produced system to equilibrate with higher multiplicities. This goes inline with the expected multi-partonic interactions, which increase for higher multiplicities in pp-collisions and are thus responsible for bringing the system towards thermodynamic equilibrium [84]. It is discovered that T_{kin} follows a mass dependent pattern and acquires higher values for heavier particles. This goes inline with the fact that heavier particles freezeout earlier in time. However, we notice multi-

plicity independent behaviour of T_{kin} particularly for lighter hadrons. The radial flow (β) parameter is also extracted in this study, which is observed higher for lighter particles. This observation reveals hydrodynamic behaviour of particles. The near-multiplicity independent behaviour of radial flow velocity, β is an important observation and this needs further investigations. Inclusion of radial flow in Taylor series expansion of Tsallis distribution function in power of $(q - 1)$, gives an analytical way in describing the transverse momentum spectra up to 3 GeV/c Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV. In addition, it could be useful in analyses where a comparison and detailed investigation is needed when comparing the Tsallis and the Boltzmann distributions. We have also studied the speed of sound as a function of temperature using the Tsallis non-extensive statistics to investigate more about the type of phase transition near the Hagedorn temperature. For small values of q and large mass cut-off, the speed of sound at first increases as a function of temperature, reaches a maximum value and then decreases, which is indicative of a phase transition. In a first order phase transition scenario, it is expected that the speed of sound drops to zero at the phase transition point. This behaviour, however, vanishes for higher values of $q \sim 1.13$. All these studies show the extensive application of the Tsallis non-extensive statistics and BGBW in high energy physics. In addition, we study the isothermal compressibility for hot and dense hadron gas using HRG and EV-HRG models. Isothermal compressibility decreases with the increasing temperature which suggest that the system at higher temperature behaves as incompressible matter. It is also observed that isothermal compressibility becomes lower for higher mass cut-offs and larger baryochemical potential. This again emphasises that, it is difficult to compress the matter with the addition of more number of baryons and heavier resonances. The effect of baryochemical potential is more prominent at lower temperatures. A collision energy dependence of κ_T is studied in EV-HRG, HRG and HRG+HS models. Isothermal compressibility initially decreases rapidly with collision energy and it becomes saturated beyond $\sqrt{s_{\text{NN}}} = 10\text{-}20$ GeV. This suggests a very different nature of the system formed in lower collision energies. Significant differences in the results from EV-HRG, HRG and HRG+HS are found for $\sqrt{s_{\text{NN}}} > 10$ GeV. We emphasise the need to have more

experimental data as a function of collision energy from very low to high energies. Addition of Hagedorn mass spectrum in the HRG model reduces the isothermal compressibility and hence it approaches closer to experimental data at higher collision energies in comparison to the other theoretical calculations. This signifies the use of Hagedorn mass spectrum into the system while studying the isothermal compressibility. Further, we have explored the application of non-extensive statistics in the field of condensed matter physics. The deviation of specific heat in the free electron model from the experimentally observed value in a number of metals is explained using the Tsallis form of FD distribution function. With the exception of Ag, in all cases considered, the non-extensive parameter, q , has been found to be greater than unity. This is mainly due to long range correlations present in the electron system. Local temperature fluctuations could also be contributing to this. This is the reason why a deviation of q from unity can explain a higher effective mass of the electrons compared to the free electron value. We have also obtained a nice systematic dependence of m^*/m as a function of the NE parameter q , a behaviour similar to which, has not been reported.

Chapter 5

Summary and Outlook

In this thesis, we started with a motivation to study pp collisions at the LHC energies in contrast to the QGP like behaviour in heavy-ion collisions. A brief discussion on the QCD phase diagram and possible phase transition and locate of CP is discussed. Few of the QGP signatures are discussed in brief. pp-collisions which usually serves as a baseline to study a medium formation in heavy-ion collisions is discussed. Then we focus on the production of resonances in pp collisions, namely K^{*0} and ϕ . In Chapter 2, we have discussed about the ALICE detectors used for our data analysis with more focus to ITS, TPC, TOF and V0 detectors. In Chapter 3, we present the study of K^{*0} and ϕ mesons production in pp-collisions with the ALICE detector at the LHC. The application of Tsallis non-extensive statistics and BGBW model in high-energy physics is provided in Chapter 4, along with its potential application in condensed matter physics as well. The important findings of this thesis are summerized below:

- The measurements of K^{*0} and ϕ production at mid-rapidity in minimum bias pp collisions at $\sqrt{s} = 8$ TeV in the range $0 < p_T < 20$ GeV/ c for K^{*0} and $0.4 < p_T < 16$ GeV/ c for ϕ has been presented.
- In comparison to other LHC energies, hardening of the p_T spectra is observed with increase in collision energy.
- The K^{*0}/π and ϕ/π ratios seem to be independent of collision energy within systematic uncertainties. This indicates that there is no strangeness enhancement in minimum bias pp collisions as the collision energy is increased.

- Similar behaviour is observed for the K^{*0}/K and ϕ/K ratios as a function of collision energy. Also, no significant energy dependence of the ϕ/K^{*0} ratio in minimum bias pp collisions at LHC energies is observed, which suggests that there is no strong energy dependence of the chemistry of the system. Finally, the p-Pb measurements at $\sqrt{s_{NN}} = 8.16$ TeV are in preparation and these pp results will serve as baseline for those in p-Pb collisions.
- The Tsallis distribution provides a complete description of identified particle spectra produced in pp -collisions at $\sqrt{s} = 7$ TeV upto very high- p_T .
- The variable T , shows a systematic increase with multiplicity, the heaviest baryons showing the steepest increase. This is an indication of a mass hierarchy in particle freeze-out.
- The notable variation of the non-extensive parameter, q which decreases towards the value 1 as the multiplicity increases. This effect is more significant for heavy mass particles. This shows the tendency of the produced system to equilibrate with higher multiplicities. This goes inline with the expected multi-partonic interactions, which increase for higher multiplicities in pp-collisions and are thus responsible for bringing the system towards thermodynamic equilibrium [84].
- Conclusively, we find that BWBG explains the transverse momentum spectra upto $\simeq 3$ GeV/ c for high multiplicity pp -collisions with an appreciable χ^2/ndf , while the Tsallis statistics describe the complete spectra for all the event classes.
- Inclusion of radial flow in Taylor series expansion of Tsallis distribution function in power of $(q - 1)$, gives an analytical way in describing the transverse momentum spectra up to 3 GeV/ c Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. In addition, it could be useful in analyses where a comparison and detailed investigation is needed when comparing the Tsallis and the Boltzmann distributions.
- The speed of sound (c_s) in a hadronic medium using non-extensive statistics has been calculated for a hadron resonance gas, as the Tsallis statistics

allows for the exploration of systems which are away from thermal equilibrium.

- For higher q -values in non-extensive statistics, c_s^2 increases near the Hagedorn's limiting temperature as compared to the extensive Boltzmann statistics.
- The phase transition in non-extensive systems is achieved at lesser temperatures as compared to the systems following extensive statistics.
- Our results leave open the possibility that there exists a special value of q where a phase transition is no longer present. Taking a mass cut-off of 2.5 GeV in the physical resonance gas, we have studied the c_s^2 as a function of system temperature, for different q -values and found that for q -values higher than 1.13, all criticality disappears. The present study bears significant implications at the LHC and higher energies, where one expects to study event-by-event temperature fluctuation specifically for the systems which stay away from thermal equilibrium.
- We have estimated the isothermal compressibility for hot and dense hadron gas using HRG and EV-HRG models. Here, hadrons and their resonances of masses only upto 2 GeV are considered in the system. We study the mass as well as baryon chemical potential dependence of κ_T as a function of temperature. We have also shown the effect of Hagedorn mass spectrum on κ_T after including the continuous mass spectra in the grand canonical partition function for hadron gas in HRG model. The lower limit of mass cut-off is set to 2 GeV in the mass integration of Hagedorn mass spectrum while the upper mass cut-offs vary from 3 – 10 GeV.
- Isothermal compressibility decreases with the increasing temperature which suggest that the system at higher temperature behaves as incompressible matter.
- It is also observed that isothermal compressibility becomes lower for higher mass cut-offs and larger baryochemical potential. This again emphasises that, it is difficult to compress the matter with the addition of more number

of baryons and heavier resonances. The effect of baryochemical potential is more prominent at lower temperatures.

- The effect of continuous mass spectra is seen on κ_T , where heavier resonances having masses above 2 GeV are included into the system. Addition of more number of resonances lowers the values of isothermal compressibility.
- A collision energy dependence of κ_T is studied in EV-HRG, HRG and HRG+HS models. Isothermal compressibility initially decreases rapidly with collision energy and it becomes saturated beyond $\sqrt{s_{NN}} = 10\text{-}20$ GeV. This suggests a very different nature of the system formed in lower collision energies. Significant differences in the results from EV-HRG, HRG and HRG+HS are found for $\sqrt{s_{NN}} > 10$ GeV. We emphasise the need to have more experimental data as a function of collision energy from very low to high energies.
- Addition of Hagedorn mass spectrum in the HRG model reduces the isothermal compressibility and hence it approaches closer to experimental data at higher collision energies in comparison to the other theoretical calculations. This signifies the use of Hagedorn mass spectrum into the system while studying the isothermal compressibility.
- We observe that the produced hadronic matter is more compressible at FAIR and NICA collision energies as compared to RHIC and LHC energies.
- We have explored the application of non-extensive statistics in the field of condensed matter physics. The deviation of specific heat in the free electron model from the experimentally observed value in a number of metals is explained using the Tsallis form of FD distribution function.
- With the exception of Ag, in all cases considered, the non-extensive parameter, q , has been found to be greater than unity. This is mainly due to long range correlations present in the electron system. Local temperature fluctuations could also be contributing to this. This is the reason why a deviation of q from unity can explain a higher effective mass of the electrons compared

to the free electron value. We have also obtained a nice systematic dependence of m^*/m as a function of the NE parameter q , a behaviour similar to which, has not been reported.

The pp collisions have been studied for a long time at the LHC, however, observables like long-range angular correlation and flow like pattern in high multiplicity pp collisions recently discovered at the LHC open up new challenges for the physicists. The origin of these effects needs proper investigation as it has been shown from hydro calculations, where a hot and dense QCD medium is implicitly assumed can describe many features of the data. However, this also can be achieved from initial state effects. Event shape studies at the LHC allows extracting more information from the data by separating the jetty events (high p_T -jets) from the isotropic events (low- Q^2 partonic scatterings). The influence of multi-partonic interactions on the final state in pp collisions needs to be understood with the study of transverse sphericity for both in data and models as well. These results can help to understand more on the jet production, identified particle ratios and the rise of mean transverse-momenta of charged particles in small systems. This suggests that more differential studies using event shapes would reveal interesting features, which could be exploited to get physics information as well as to improve models in the MC event generators.

Chapter 6

Appendix

6.1 Expansion in $(q - 1)$ of Tsallis distribution

The Taylor expansion is done efficiently by the following change of variables:

$$q - 1 \equiv x; \frac{E - \mu}{T} \equiv \Phi \quad (6.1)$$

$$1 + (q - 1) \frac{E - \mu}{T} \equiv 1 + x \Phi \equiv f(x) \quad (6.2)$$

$$-\frac{q}{q - 1} = -\frac{1 + x}{x} \equiv g(x) \quad (6.3)$$

The relevant distribution function can then be written as

$$\begin{aligned} \left[1 + (q - 1) \frac{E - \mu}{T} \right]^{-\frac{q}{q-1}} &= [1 + x\Phi]^{-\frac{1+x}{x}} \\ &= f(x)^{g(x)} \\ &= \mathcal{F}(x) \end{aligned} \quad (6.4)$$

Expanding $\mathcal{F}(x)$ in Taylor series about $x = 0$, we get

$$\mathcal{F}(x = 0) = e^{-\Phi} \quad (6.5)$$

$$\left. \frac{d\mathcal{F}(x)}{dx} \right|_{x=0} = \frac{1}{2}(-2 + \Phi)\Phi e^{-\Phi} \quad (6.6)$$

$$\left. \frac{d^2\mathcal{F}(x)}{dx^2} \right|_{x=0} = \frac{\Phi^2}{12} (24 - 20\Phi + 3\Phi^2) e^{-\Phi} \quad (6.7)$$

for $\mathcal{O}(x^0)$, $\mathcal{O}(x)$ and $\mathcal{O}(x^2)$ respectively. The final results can be summarized in

the following equation

$$\begin{aligned}
& \left[1 + (q-1) \frac{E-\mu}{T} \right]^{-\frac{q}{q-1}} \\
& \simeq e^{-\frac{E-\mu}{T}} \\
& + (q-1) \frac{1}{2} \frac{E-\mu}{T} \left(-2 + \frac{E-\mu}{T} \right) e^{-\frac{E-\mu}{T}} \\
& + \frac{(q-1)^2}{2!} \frac{1}{12} \left[\frac{E-\mu}{T} \right]^2 \\
& \times \left[24 - 20 \frac{E-\mu}{T} + 3 \left(\frac{E-\mu}{T} \right)^2 \right] e^{-\frac{E-\mu}{T}} \\
& + \dots
\end{aligned} \tag{6.8}$$

$$\tag{6.9}$$

For completeness we also quote the following result

$$\begin{aligned}
& \left[1 + (q-1) \frac{E-\mu}{T} \right]^{-\frac{1}{q-1}} \\
& \simeq e^{-\frac{E-\mu}{T}} \\
& + (q-1) \frac{1}{2} \left(\frac{E-\mu}{T} \right)^2 e^{-\frac{E-\mu}{T}} \\
& + \frac{(q-1)^2}{2!} \frac{1}{12} \left[\frac{E-\mu}{T} \right]^2 \\
& \times \left[-8 \left(\frac{E-\mu}{T} \right) + 3 \left(\frac{E-\mu}{T} \right)^2 \right] e^{-\frac{E-\mu}{T}} \\
& + \dots
\end{aligned} \tag{6.10}$$

6.2 Tsallis Thermodynamics

6.2.1 Particle number density n

Up to $\mathcal{O}(q-1)$, the number density can be written as:

$$\begin{aligned}
n = & \frac{g}{(2\pi)^3} \int d^3p \left[e^{-\frac{E-\mu}{T}} + (q-1) \frac{E-\mu}{2T} \right. \\
& \left. \left(-2 + \frac{E-\mu}{T} \right) e^{-\frac{E-\mu}{T}} \right]
\end{aligned} \tag{6.11}$$

We define,

$$\begin{aligned}
p^2 + m^2 &= E^2; \\
\frac{E}{T} &= \omega, \frac{m}{T} = a, \frac{\mu}{T} = b
\end{aligned} \tag{6.12}$$

And hence,

$$\begin{aligned}
n &= n^B + \frac{g(q-1)e^{\frac{\mu}{T}}T^3}{4\pi^2} \int d\omega \, \omega (\omega^2 - a^2)^{\frac{1}{2}} \\
&\quad (-2\omega + 2b + \omega^2 + b^2 - 2\omega b) e^{-\omega} \\
n &= n^B + \frac{g(q-1)e^{\frac{\mu}{T}}T^3}{4\pi^2} [-6a^2 K_2(a) - 2a^3 K_1(a) \\
&\quad - 4a^2 b K_2(a) + 3a^3 K_3(a) + a^4 K_2(a) + a^2 b^2 K_2(a) \\
&\quad - 2a^3 b K_1(a)]
\end{aligned} \tag{6.13}$$

Here the following form of the modified Bessel function of second kind [?] is used:

$$K_n(a) = \frac{2^{n-1}(n-1)!}{(2n-2)!a^n} \int_a^\infty d\omega \, \omega (\omega^2 - a^2)^{n-\frac{3}{2}} e^{-\omega} \tag{6.14}$$

On the other hand, by using Boltzmann Statistics one obtains,

$$n^B = \frac{ge^{\frac{\mu}{T}}T^3 a^2 K_2(a)}{2\pi^2} \tag{6.15}$$

6.2.2 Pressure

To first order in $(q-1)$ the pressure is given by:

$$\begin{aligned}
P &= \frac{g}{(2\pi)^3} \int d^3p \, \frac{p^2}{3E} \left[e^{-\frac{E-\mu}{T}} + (q-1) \frac{E-\mu}{2T} \right. \\
&\quad \left. \left(-2 + \frac{E-\mu}{T} \right) e^{-\frac{E-\mu}{T}} \right]
\end{aligned} \tag{6.16}$$

Using the same definitions as in Eq. 6.12, we get

$$\begin{aligned}
P &= P^B + \frac{g(q-1)e^{\frac{\mu}{T}}T^4}{12\pi^2} \int d\omega (\omega^2 - a^2)^{\frac{3}{2}} \\
&\quad (-2\omega + 2b + \omega^2 + b^2 - 2\omega b) e^{-\omega} \\
&= P^B + \frac{g(q-1)e^{\frac{\mu}{T}}T^4}{4\pi^2} [a^4 K_2(a) + 3a^3 K_3(a) \\
&\quad - 2a^3 b K_3(a) + a^2 b^2 K_2(a) + 2a^2 b K_2(a)]
\end{aligned} \tag{6.17}$$

Using Eq. 6.12 along with Eq. 6.14 as well as another representation of the modified Bessel function,

$$K_n(a) = \frac{2^n n!}{(2n)! a^n} \int_a^\infty d\omega (\omega^2 - a^2)^{n-\frac{1}{2}} e^{-\omega} \tag{6.18}$$

Eq. 6.14 can be obtained from Eq. 6.18 by dint of partial integration. The Boltzmann pressure density is, similarly given by:

$$P^B = \frac{g e^{\frac{\mu}{T}} T^4 a^2 K_2(a)}{2\pi^2} \tag{6.19}$$

6.2.3 Energy density ϵ

$$\begin{aligned}
\epsilon &= \frac{g}{(2\pi)^3} \int d^3p E \left[e^{-\frac{E-\mu}{T}} + (q-1) \frac{E-\mu}{2T} \right. \\
&\quad \left. \left(-2 + \frac{E-\mu}{T} \right) e^{-\frac{E-\mu}{T}} \right]
\end{aligned} \tag{6.20}$$

Using Eq. 6.12 along with Eq. 6.14,

$$\begin{aligned}
\epsilon &= \epsilon^B + \frac{g(q-1)e^{\frac{\mu}{T}}T^4}{4\pi^2} [9a^3 K_3(a) + 4a^4 K_2(a) + a^5 K_1(a) \\
&\quad + 2b (3a^2 K_2(a) + a^3 K_1(a) - 3a^3 K_3(a) + a^4 K_2(a)) \\
&\quad b^2 (3a^2 K_2(a) + a^3 K_1(a))]
\end{aligned} \tag{6.21}$$

with

$$\epsilon^B = \frac{g e^{\frac{\mu}{T}} T^4}{2\pi^2} (3a^2 K_2(a) + a^3 K_1(a)) \tag{6.22}$$

6.2.4 Checking Thermodynamic Consistency

We know that for thermodynamic variables n and P at non-zero μ we have,

$$n = \partial P / \partial \mu \quad (6.23)$$

$$\begin{aligned} n^B + (q-1)n^1 + \dots &= \partial P^B / \partial \mu + \\ & (q-1)\partial P^1 / \partial \mu + \dots \end{aligned} \quad (6.24)$$

where the superscripts denote the order of $(q-1)$ in the expansion. Since the terms in the expansion are linearly independent, the identity given by Eq. 6.23 is to be satisfied at every order of $(q-1)$. Since our expansion is up to $\mathcal{O}(q-1)$, we will check the identity for that order now. Rearranging and reordering the $\mathcal{O}(q-1)$ terms in Eq. 6.13 we get,

$$\begin{aligned} n^1 &= \frac{ge^{\frac{\mu}{T}}T^3}{4\pi^2} [a^3K_3(a) + a^4K_2(a) \\ & - 4a^2bK_2(a) + 2a^3b(K_3(a) - K_1(a)) \\ & - 2a^3bK_3(a) + a^2b^2K_2(a) - 6a^2K_2(a) \\ & + 2a^3(K_3(a) - K_1(a))] \end{aligned} \quad (6.25)$$

Using the recursion relation for the modified Bessel's functions,

$$K_{n+1}(a) - K_{n-1}(a) = \frac{2n}{a}K_n(a) \quad (6.26)$$

$$\begin{aligned} n^1 &= \frac{g(q-1)e^{\frac{\mu}{T}}T^3}{4\pi^2} [a^3K_3(a) + a^4K_2(a) \\ & + 4a^2bK_2(a) - 2a^3bK_3(a) + a^2b^2K_2(a) \\ & + 2a^2K_2(a)] \\ &= \frac{\partial P^1}{\partial \mu} \end{aligned} \quad (6.27)$$

Hence, proved.

6.3 Momentum Distribution

The invariant particle yield is given by,

$$E \frac{dN}{d^3p} = \mathcal{C} E \left[1 + (q-1) \frac{E-\mu}{T} \right]^{-\frac{q}{q-1}} \quad (6.28)$$

Where $\mathcal{C} = \frac{gV}{(2\pi)^3}$. Assuming $q-1 \ll 1$ we are to expand it in Taylor series. Let,

$$\begin{aligned} \frac{dN}{p_T dp_T dy d\phi} = \mathcal{C} E & \left[e^{-\Phi} + \frac{x}{1!} \frac{1}{2} (-2 + \Phi) \Phi e^{-\Phi} \right. \\ & \left. + \frac{x^2}{2!} \frac{\Phi^2}{12} (24 - 20\Phi + 3\Phi^2) e^{-\Phi} + \dots \right] \end{aligned} \quad (6.29)$$

where y is rapidity and ϕ is the azimuthal angle of emission. Now, parameterizing energy E in terms of rapidity y i.e. putting $E = m_T \cosh y$ and integrating over y we get Eq. 4.28.

6.4 Flow in Tsallis Distribution

We use the following ansatz (in cylindrical polar coordinates) for introducing flow inside our calculations:

$$p^\mu = (m_T \cosh y, p_T \cos \phi, p_T \sin \phi, m_T \sinh y) \quad (6.30)$$

$$u^\mu = (\gamma \cosh \zeta, \gamma v \cos \alpha, \gamma v \sin \alpha, \gamma \sinh \zeta) \quad (6.31)$$

where $(\zeta)y$ is the (space-time)rapidity of particles (fluid-element) and v is the velocity of fluid. Now, to include flow inside the Tsallis distribution, we replace $E \rightarrow p^\mu u_\mu$ assuming temperature to be scalar. The dot product then becomes,

$$p^\mu u_\mu = \gamma m_T \cosh(y - \zeta) - \gamma v p_T \cos(\phi - \alpha) \quad (6.32)$$

Now putting Eq. 6.32 in Eq. 4.27 up to $\mathcal{O}(q-1)$ and integrating over ϕ

and ζ we get,

$$\begin{aligned}
\frac{1}{p_T} \frac{dN}{dp_T dy} &= \frac{gV}{(2\pi)^2} \\
&\left\{ 2T[rI_0(s)K_1(r) - sI_1(s)K_0(r)] \right. \\
&\quad - (q-1)Tr^2I_0(s)[K_0(r) + K_2(r)] \\
&\quad + 4(q-1)TrsI_1(s)K_1(r) \\
&\quad - (q-1)Ts^2K_0(r)[I_0(s) + I_2(s)] \\
&\quad + \frac{(q-1)}{4}Tr^3I_0(s)[K_3(r) + 3K_1(r)] \\
&\quad - \frac{3(q-1)}{2}Tr^2s[K_2(r) + K_0(r)]I_1(s) \\
&\quad + \frac{3(q-1)}{2}Ts^2r[I_0(s) + I_2(s)]K_1(r) \\
&\quad \left. - \frac{(q-1)}{4}Ts^3[I_3(s) + 3I_1(s)]K_0(r) \right\} \quad (6.33)
\end{aligned}$$

where

$$r \equiv \frac{\gamma m_T}{T} \quad (6.34)$$

$$s \equiv \frac{\gamma v p_T}{T} \quad (6.35)$$

$I_n(s)$ is the modified Bessel function of first kind.

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