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Study of the production of (anti-)deuterons in pp collisions at the LHC with the ALICE experiment

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Abstract

At the LHC an abundant production of light (anti-)nuclei, i.e. with a mass number $A \leq 4$, is observed. However, although abundantly produced, their production mechanisms are still not completely understood. Indeed, light (anti-)nuclei are characterised by a binding energy of the order of 1 MeV per nucleon and this value is very low if compared with the chemical freeze-out temperature of a Pb–Pb collision ($T_{ch} \sim 150$ MeV). Therefore, in principle the observation of such fragile objects is not expected and a comparison of the experimental data with the theoretical predictions can help to shed light on the production mechanism of these loosely bound states.

ALICE is the most suitable LHC experiment for the study of light (anti-)nuclei, thanks to its excellent particle identification (PID) capabilities. In particular, the identification of light nuclei is possible via the measurement of the specific energy loss dE/dx in the Time Projection Chamber (TPC) and with the time information provided by the Time Of Flight (TOF) detector.

Among light (anti-)nuclei, (anti-)deuterons are the lightest and consequently the most abundantly produced. The ALICE Collaboration has measured the production spectra of (anti-)deuterons as a function of charged particle multiplicity $dN_{ch}/d\eta$ in proton-proton (pp) collisions at the energy of $\sqrt{s} = 7$ TeV, in proton-lead (p–Pb) collisions at $\sqrt{s_{NN}} = 5.02$ TeV and in lead-lead (Pb–Pb) collisions at $\sqrt{s_{NN}} = 2.76$ TeV.

In this work, the (anti-)deuteron production spectra in pp collisions at $\sqrt{s} = 13$ TeV as a function of the event charged particle multiplicity and in the Minimum Bias data sample is presented. Thanks to the extremely large data sample collected in 2016 and 2017 by ALICE, consisting of almost 1 billion pp collisions, it has been possible to study the (anti-)deuteron production in very fine multiplicity classes and in particular reaching multiplicities that are similar to those observed in p–Pb collisions. This is a crucial element to have a comparison of different colliding systems and to understand if the production mechanism of light (anti-)nuclei is similar in small and medium-large colliding systems.

Relying on the measurement of the (anti-)proton production spectra, the coalescence parameter B_2 , which is related to the probability to form a deuteron via coalescence, and the ratio between p_T -integrated yields of (anti-)deuterons and (anti-)protons d/p are computed. The measurement of B_2 and d/p as a function of charged particle multiplicity is a useful tool for shedding light on the dependence of the production mechanisms on the system size. The results are also relevant for background studies in the search for dark matter via the measurement of (anti-)nuclei in space and as input for the understanding of the formation of QCD bound states in high energy hadron physics.

The results as a function of the charged particle multiplicity are compared with the predictions of the two available classes of phenomenological models, namely the statistical hadronisation model and the coalescence model. This is the first deuteron analysis within the ALICE Collaboration in which such comparison is performed. The results have been approved by the ALICE Collaboration and have been submitted to publication [1].

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Probing strongly-interacting matter at the LHC

Strong interaction is one of the four fundamental interactions known so far. It is responsible for keeping together protons and neutrons (nucleons) inside nuclei, and quarks, the fundamental constituents of matter, inside nucleons. The study of strongly-interacting matter is crucial to understand the evolution of the Universe in its first stages after the Big Bang. The characteristics of strong interaction, which will be described in this section, are very peculiar: in a high energy-density environment, such as the early Universe [2–5], quarks cannot bind together. However, when the Universe expands the energy-density decreases and a transition from this phase to ordinary matter takes place. The signs of this transition cannot be observed in the current Universe and therefore such conditions must be recreated in the laboratory [6].

1.1 Quantum Chromodynamics

Quantum Chromodynamics (QCD) is the gauge field theory that describes strong interaction within the Standard Model [7]. The name chromodynamics derives from *colour*, which is the name of the QCD charge.

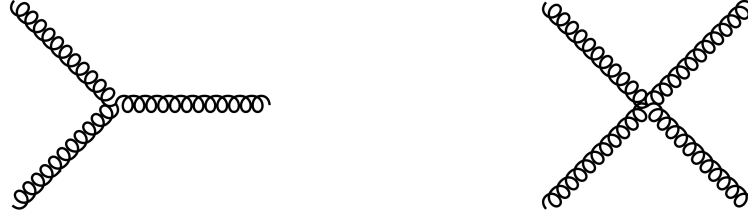
The QCD Lagrangian has an explicit SU(3) symmetry, which results in eight massless mediator bosons (gluons) [8]. It has the form:

$$\mathcal{L}_{QCD} = -\frac{1}{4}G_{\mu\nu}^a G_a^{\mu\nu} + \sum_f \bar{\Psi}_f (i\gamma^\mu \partial_\mu + g_s \frac{\lambda_a}{2} \gamma^\mu A_\mu^a - m_f) \Psi_f \quad (1.1)$$

where g_s is the coupling constant of strong interaction, a is the group index, A_μ^a is the gluon field associated to the SU(3) generator λ^a and f is the quark flavour. It is possible to see how the flavour symmetry is explicitly broken by the quark mass.

The main difference between QCD and Quantum Electrodynamics (QED), the gauge field theory that describes the electromagnetic interaction, arises from the non-Abelian nature of the symmetry group SU(3). The gluon tensor has the form:

$$G_a^{\mu\nu} = \partial^\mu A_a^\nu - \partial^\nu A_a^\mu - g_s f_{abc} A^{b\mu} A^{c\nu}. \quad (1.2)$$



$$(a) -\frac{g_s}{2} f_{abc} (\partial^\mu A^{\nu a} - \partial^\nu A^{\mu a}) A_\mu^b A_\nu^c$$

$$(b) -\frac{g_s^2}{4} f_{abe} f_{cde} A_{a\mu} A_{b\nu} A_c^\mu A_d^\nu$$

Fig. 1.1.: Three and four-gluon interaction vertices.

where f_{abc} is a structure constant of SU(3). While in QED the unique structure-constant of U(1), which is Abelian, is null, f_{abc} is not null and the additional term in the tensor brings two self-interaction vertices for gluons, as shown in Fig. 1.1.

The gluon self-interaction vertices bring loop corrections to the gluon propagator, leading to the so-called *anti-screening*: the intensity of strong interaction increases with distance. The reason behind this effect is that gluons are not colour-singlets, but they carry colour-charge. In QED, instead, photons are neutral objects and loop-corrections to their propagator have the opposite effect.

The coupling constant of the strong interaction $\alpha_s = g_s^2/4\pi$ evolves with distance or, inversely, with transferred momentum Q [9, 10], according to equation 1.3:

$$\alpha_s(Q^2) = \frac{\alpha_s(\mu^2)}{1 + \alpha_s(\mu^2) \frac{33-2n_f}{12\pi} \ln \frac{Q^2}{\mu^2}} \quad (1.3)$$

where n_f is the number of flavours and μ is the renormalisation scale of the theory. A summary of the measurements of α_s as a function of the energy scale Q is shown in Fig. 1.2.

Due to the transferred-momentum dependence of α_s , two different regimes can be observed. For high Q^2 , α_s goes to zero and QCD becomes a free theory. This regime is called *asymptotic freedom*. On the other hand, for low Q^2 the coupling constant α_s assumes high values and, as a consequence, quarks are bound in hadrons. This regime is called *confinement*.

For $Q^2 \gg \mu^2$, $\alpha_s \ll 1$ and a perturbative approach can be used. With perturbative QCD (pQCD) it is possible to calculate transition elements of the scattering matrix. However, at low Q^2 a different approach is needed. One solution consists of evaluating the Green's functions of the QCD Lagrangian on a space-time lattice with spacing a . This method is called lattice QCD (LQCD) [11]. Via extrapolation to

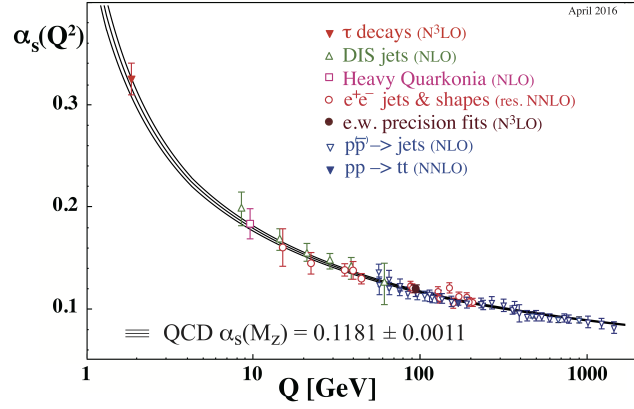


Fig. 1.2.: Summary of measurements of α_s as a function of the energy scale Q . The respective degree of QCD perturbation theory used in the extraction of α_s is indicated in brackets (NLO: next-to-leading order; NNLO: next-to-next-to leading order; res. NNLO: NNLO matched with re-summed next-to-leading logs; N3LO: next-to-NNLO). Figure taken from Ref. [4].

the continuum ($a \rightarrow 0$) it is possible to obtain results that can be compared with experimental measurements. One of the most impressive successes of LQCD in the description of soft processes of the colour interaction is the determination of the proton mass with a precision of 2% [12].

1.2 Production of matter in the early Universe

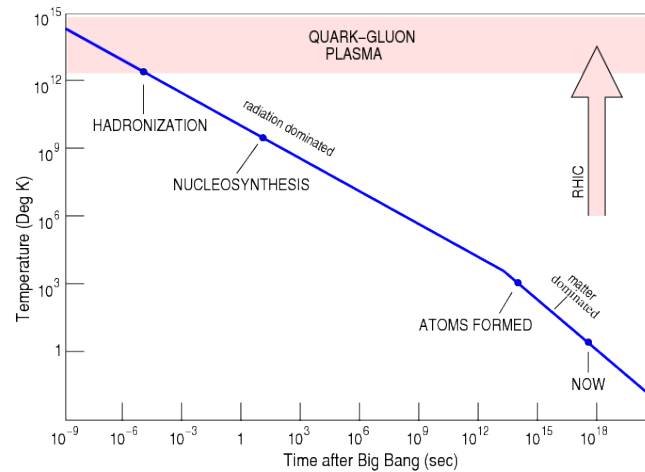


Fig. 1.3.: Temperature of the Universe as a function of time after Big Bang. The arrow on the right indicates the temperature reached in heavy-ion collision experiments.

According to the Big Bang model, the Universe expanded from an extremely small region characterised by infinite energy density and temperature [13]. In the first few

microseconds of the Universe life, the temperature was so high that quarks could not bind together to form hadrons. Instead, quarks and gluons could freely move in a phase of strongly-interacting plasma, called *Quark-Gluon Plasma* (QGP) [14].

As an effect of the expansion, the temperature and the energy density of the Universe decreased. Thus when the temperature went below the critical value $T_C \sim 150$ MeV [15], as an effect of α_s energy dependence, a phase transition took place: quarks and gluons were confined into hadrons, which are colour-singlet states. This phase transition is called *hadronisation*.

The high temperature of the Universe allowed for the interconversion of neutrons and protons via electroweak interaction: $n + e^+ \leftrightarrow \bar{\nu}_e + p$, $n + \nu_e \leftrightarrow e^- + p$ and $n \leftrightarrow e^- + \bar{\nu}_e + p$. The neutron/proton ratio was at equilibrium as long as the interconversion rate was faster than the expansion of the Universe. At $T_C \sim 1$ MeV the equilibrium of the reaction broke and n/p froze out, thereafter decreasing slowly due to the neutron decay. Although the freeze-out occurred at a temperature that was below the deuteron binding energy $E_B = 2.2$ MeV, nuclei could not still be formed for the presence of high-energy photons: the photo-destruction rate was larger than the production rate. When the temperature dropped below 100 keV, neutrons and protons could bind and form light nuclei, mainly hydrogen and helium nuclei: this process is known as *primordial nucleosynthesis* [16].

At this stage the Universe was still ionised and hence opaque to electromagnetic radiation. However, when the Universe reached a temperature of 3000K, atoms started forming and the electromagnetic radiation decoupled with a black-body spectrum. As a result of the Universe expansion, the electromagnetic radiation is red-shifted and today the wavelength is in the region of microwaves, with a temperature of 2.7 K, and constitutes the so-called *Cosmic Microwave Background* (CMB) [17].

The anisotropies in matter distribution at the time of electromagnetic-radiation decoupling kept growing under the effect of gravity and all the structures present in the current Universe, such as stars and galaxies, started forming. Evidences of the initial anisotropy in the matter distribution can be found in the CMB temperature anisotropies [18].

This brief summary of the evolution of the Universe from Big Bang, together with its temperature as a function of time, is shown in Fig. 1.3 [19]. As already mentioned above, before the decoupling of radiation the Universe was opaque and therefore everything that happened before that time can not be observed in the current Universe. In particular, the signatures of the transition between QGP and hadronic phase cannot be studied. For this reason, those conditions must be recreated in laboratory, using collider experiments.

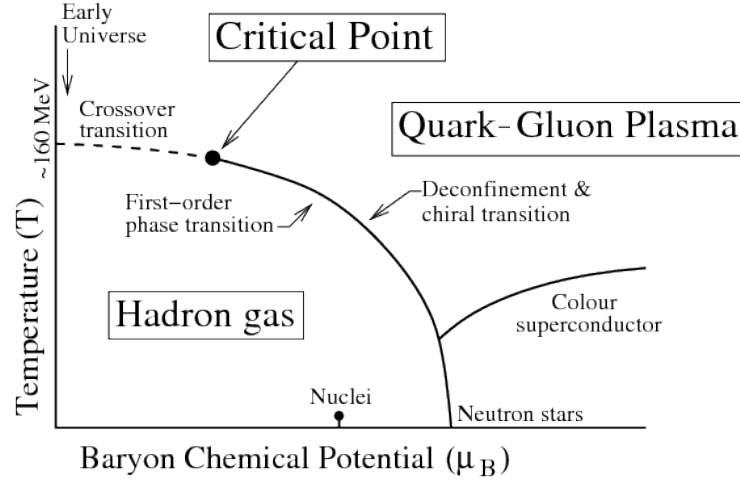


Fig. 1.4.: QCD phase diagram. Figure taken from Ref. [19].

1.3 QCD phase diagram

As a consequence of the running α_s , the properties of strongly-interacting matter are highly dependent on the mean transferred momentum in the interactions inside the medium. To describe hadronic systems, i.e. composed of quark and gluons, two thermodynamical variables are needed. The temperature T and the baryon-chemical potential μ_B , i.e. the chemical potential associated to the baryon number, are typically used.

Fig. 1.4 shows the phase diagram for QCD matter. The origin of the phase diagram corresponds to the QCD vacuum. Ordinary nuclear matter is located around $T = 0$ and $\mu_B \approx 1 \text{ GeV}$. Moving along the μ_B axis, theoretical models predict a phase transition to a deconfined state. This state has been hypothesised to be present in the core of neutron stars [20]. For $\mu_B \lesssim 1 \text{ GeV}$ we have the hadronic phase [21]. Moving along the T axis, in the region $\mu_B \approx 0$ and $T \gg \Lambda_{QCD}$, being Λ_{QCD} the QCD scale, the conditions become similar to those existing in the primordial Universe and a transition to a state of deconfined strongly interacting matter is foreseen. Indeed, as the temperature increases also the mean transferred momentum increases and the strong interaction becomes less and less intense. The transition from confined to deconfined strongly-interacting matter is expected to be a first order transition at high μ_B . On the other hand, for $\mu_B \approx 0$ a crossover transition is expected, at the temperature $T_C \sim 150 \text{ MeV}$ [15].

In the early Universe the transition happened from the QGP phase to hadronic matter at $\mu_B \sim 0.33 \text{ eV}$ [22], caused by the expansion and the consequent cooling. In order to recreate a QGP phase it is necessary to take the opposite path, i.e. to create a system characterised by an extreme energy density. A simple estimation of the

energy density for the phase transition can be obtained by considering quarks and gluons as a gas of massless particles confined in a volume V [23]. According to statistical mechanics, for a gas of massless particles the energy density ϵ is related to the temperature T by the equation:

$$\epsilon = \frac{\pi^2}{30} n_{DOF} T^4 \quad (1.4)$$

where n_{DOF} is the number of degrees of freedom of the system, which for a system of gluons and quarks is $n_{DOF} = \left(16 + \frac{21}{2} n_f\right)$, being n_f the number of quark flavours. Considering $n_f = 2$ because in ordinary matter there are only u and d quarks, $n_{DOF} = 37$. It means that for a critical temperature $T_C = 150$ MeV the energy density is $\epsilon \sim 1$ GeV/fm³.

1.4 Heavy ion collisions

The first experiments on heavy-ion collisions (HIC) were performed with a fixed target configuration at the *Alternate Gradient Synchrotron* (AGS) [24] at the Brookhaven National Laboratory and at the *Super Proton Synchrotron* (SPS) [25] at CERN, with centre-of-mass (CM) energies in the range between 2 A GeV and 18 A GeV, being A the mass number of the nuclei.

The next generation of experiments was based on ion-beam collisions. The two hadron colliders that have an active HIC program are the *Relativistic Heavy Ion Collider* (RHIC) at the Brookhaven National Laboratories (BNL) and the *Large Hadron Collider* (LHC) at CERN. RHIC and LHC can provide heavy-ion collisions with a CM energy up to 200 A GeV and 5.02 A TeV respectively. Recent reviews of the results obtained at RHIC and at LHC can be found in Ref. [26] and in Ref. [27] respectively. At the LHC an energy density of 15 GeV/fm³ was measured for the initial stage of Pb-Pb collision at the CM energy of 2.76 A TeV [27], which is well above the critical energy density of 1 GeV/fm³.

1.4.1 Collision geometry

A phenomenological description of nucleus-nucleus collisions can be given with the Glauber model [28]. Starting from the geometrical configuration of the colliding nuclei, the Glauber model describes nucleus-nucleus collisions in terms of the collisions of the constituent nucleons. The key parameter of the model is the impact parameter b , which is defined as the distance between the centres of the colliding nuclei in the transverse plane to their momenta (see Fig. 1.5).

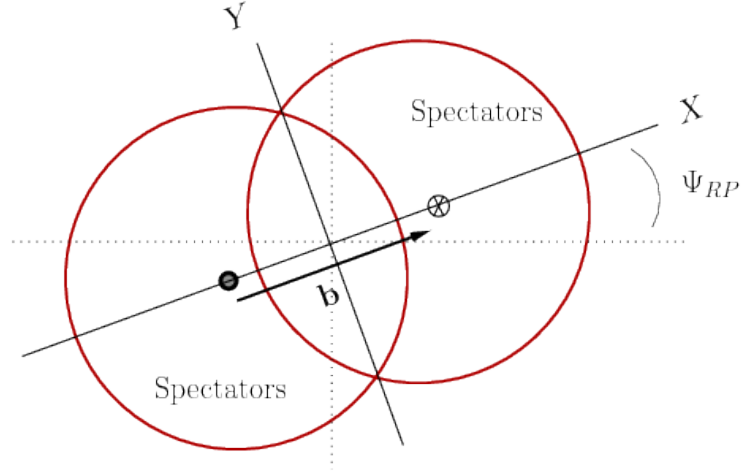


Fig. 1.5.: Definition of the impact parameter b in a heavy-ion collision. The colliding nuclei are shown in the transverse plane, which is perpendicular to their directions. The directions of the colliding nuclei are represented by two arrows in their centres, one entering and one exiting the transverse plane, respectively. Figure taken from Ref. [29].

The impact parameter b is used to define the centrality of a collision. Central collisions are characterised by a small impact parameter. In central collisions many nucleons are involved in the interaction and hence there are many binary collisions between nucleons. Moreover, central collisions are characterised by a large interaction volume and by the production of many particles in the final state. On the other hand, peripheral collisions are those with a large impact parameter and their characteristics are just the opposite of central-collision ones.

In a collision the nucleons that have at least one interaction with one of the nucleons of the other nucleus are called *participants*, while all the nucleons that do not interact are called *spectators*. Spectators proceed with little perturbation along their original direction and participants are left behind in an excited almond-shape region. Using the theory of probability, based on the nucleon-nucleon cross sections and on the density distribution of nucleons inside the nucleus, the Glauber model provides quantitative expressions to calculate the number of elementary nucleon-nucleon collisions (N_{coll}), the number of participants (N_{part}), the interaction probability and the size of the overlap region between the colliding nuclei.

1.5 Space-time evolution of the collision

When two ultra-relativistic nuclei collide, a long-lived strongly interacting system is created. One of the subjects of investigation of HIC experiments is the evolution of such a system and its properties. The current view on the evolution of HIC [30] is shown in Fig. 1.6.

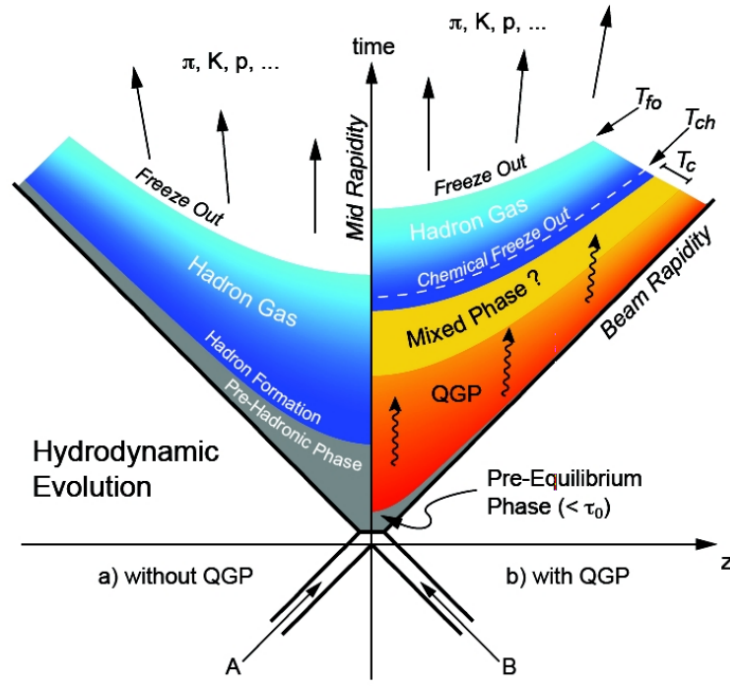


Fig. 1.6.: Evolution of a central heavy ion collision in a space-time diagram. The z direction is parallel to the beam line. The two scenarios with (right) and without (left) QGP are pointed out.

The collision can be divided into several steps:

1. **Pre-collision stage.** Before the collision, nuclei travel in the beam line. At the energy reached in the accelerator nuclei are strongly Lorentz-contracted in the laboratory reference frame (by a factor 100 and 2700 at RHIC and at LHC respectively). They can be seen as two disks.
2. **Collision.** The collision takes place: its geometry can be described as in section 1.4.1.
3. **Pre-equilibrium stage.** In the first instants of the collision hard processes takes place. They are processes characterised by a high transferred momentum between colliding partons. All the particles with high energy (i.e. high momentum and/or mass) are produced at this stage.
4. **Thermal equilibrium.** In high energy collisions, the momentum of the nuclei is so high that the constituent partons interact several times, losing energy in the mid-rapidity¹ region ($y \approx 0$) and escaping at forward rapidity ($|y| \gg 0$). Thus the system results in a hot interacting medium at mid-rapidity with vanishing baryon-chemical potential ($\mu_B \approx 0$). The baryon-chemical potential

¹rapidity is defined in section 1.6

carried out by the colliding nuclei is brought at forward rapidity by constituent quarks and spectator nucleons. At RHIC and at LHC a transition to a QGP state is expected due to the high energy-density (see Fig. 1.6 right). After some strong parton scattering, the QGP droplet reaches the equilibrium.

5. **Phase transition.** The QGP droplets expands under the effects of the thermal pressure gradients at the system boundaries. This phase of expansion of the QGP is modelled with relativistic hydrodynamics [31]. As a consequence of the expansion, the system cools down, crossing the phase boundary between deconfined and confined strongly interacting matter. When the critical temperature T_c between the two phases is reached, the hadronisation takes place and the system gradually evolves into an interacting *hadron resonance gas*.
6. **Chemical freeze-out.** As the system expands and cools down, both inelastic and elastic interactions among hadrons keep taking place. The instant when the transferred momentum between hadrons is not sufficient for inelastic interaction is called *chemical freeze-out*. At the chemical freeze-out the abundances of the particle species are fixed.
7. **Kinetic freeze-out.** The system keeps expanding and cooling down and, at the *kinetic freeze-out*, hadrons stop interacting. At the chemical freeze-out the kinematics of the particles is fixed.
8. **Late stage.** Hadrons escape the interaction region without further interactions.

In case the energy-density is not sufficient to create a QGP droplet, a hadron gas is directly formed and the system evolves in the same way as from point 6 (see Fig. 1.6 left). For a pp collision, the evolution of the system is similar to a Pb–Pb collision, but the lower production of quarks and gluons in the initial stages determines a qualitatively different final state.

1.5.1 For each stage an appropriate probe

Experimentally, the observables can be measured only at the state in which they are after the last step of the aforementioned system evolution. However, it is possible to have an indirect look into intermediate stages of the evolution through the comparison of final-state observables with the predictions of theoretical models. The observables can be divided in three macro-categories: hard probes, electroweak probes and soft probes.

Hard probes are produced in the first stages of the collisions, when processes with high transferred momentum occur: either quarks and gluons with high momentum or heavy-flavoured quarks (charm and beauty) are created. For the same reason, their production rates are computable with pQCD. Since with the evolution of the system the energy density decreases, hard probes can be produced only in the first stages of the collision and they live through the successive stages. Therefore the study of hard probes allows us to inquire into the mechanisms driving the parton propagation and the energy loss in the QGP. The hadronisation of quarks and gluons produced in the early stages of the collision brings to the formation of open heavy-flavoured particles, quarkonia and jets. Open heavy-flavoured particles are hadrons composed of a heavy quark (c or b) and light quarks (d, u or s), such as D and B mesons. Similarly, quarkonia are hadrons formed by a quark and an anti-quark with the same flavour such as the J/Ψ ($c\bar{c}$). Jets are instead narrowly collimated bundles of hadrons. The effects of the energy loss and the dynamics of heavy-quark hadronisation can be studied with the comparison of the transverse-momentum differential production yields in nucleus-nucleus collisions with those obtained in pp collisions, in which QGP is not expected to form.

Electroweak probes include leptons, photons, Z and $W^\pm$ bosons. Electroweak probes are colour blind probes: at the extreme energy-density that characterise a HIC the dominant interaction is the strong one and all the other are negligible. For this reason, they do not interact with the medium and they carry the information about the initial stages of the collisions without modifications. Direct photons, i.e. those not produced in particle decays, are particularly interesting because they include not only photons produced in the early stages of the collisions but also thermal photons emitted by the QGP. Their study is extremely interesting because their energy is related to the temperature of the QGP. However, thermal photons are not discernible from direct photons coming from other sources and all direct photons must be studied as a whole. The first measurement of direct-photon production in Pb–Pb collisions by the ALICE collaboration can be found in [32].

Finally, soft probes are low momentum hadrons and they represent the majority of the particles produced in a collision ($\approx 99\%$). They are produced in the latest stages of the collision, when the energy density of the system is highly degraded and therefore their production rates must be calculated with non-perturbative QCD. The topic of this thesis belongs to this category.

Only the soft probes which are relevant for this analysis will be described in the following, after the definition of some important variables in section 1.6. All the other observables will not be discussed.

1.6 Some useful definitions

The observables reported in the following sections are referred to a right-handed Cartesian system, centred in the collision point, with the z axis parallel to the beam direction, the x axis aligned with the accelerator plane, pointing towards the centre of the ring, and the y axis pointing upwards, according to the chirality of the system. The z direction is referred as *longitudinal* dimension and the xy plane as the *transverse plane*. For example, the momentum $\vec{p}$ of a particle can be decomposed into its longitudinal momentum with module p_z and its transverse momentum with module $p_T = \sqrt{p_x^2 + p_y^2}$. The polar angle θ is expressed with respect to the z axis.

The rapidity of a particle is defined as:

$$y = \frac{1}{2} \ln \frac{E + p_z}{E - p_z} \quad (1.5)$$

where E is the energy of the particle and p_z is its longitudinal momentum. The pseudorapidity η is instead defined as:

$$\eta = -\ln \tan(\theta/2) = \frac{1}{2} \ln \frac{|\vec{p}| + p_z}{|\vec{p}| - p_z} \quad (1.6)$$

At high energy $E \sim |\vec{p}|$ and hence $\eta \sim y$.

1.7 Charged particle multiplicity and system size

Charged particle multiplicity dN_{ch}/dy (from here on just *multiplicity*), i.e. the number of charged particle produced per rapidity unity, is an extremely important observable. It can be used to estimate the centrality of a Pb-Pb collision. Indeed, the Glauber Model (see section 1.4.1) can provide the relation between impact parameter and multiplicity. In this way it is possible to get an indirect measurement of the impact parameter from the measurement of multiplicity. Moreover, multiplicity is defined for each collision system and hence it allows the direct comparison of observable across different collision systems.

Multiplicity can be related to the size of the colliding system. The radius of the collision system is usually estimated from the study of two-particle correlations using the Hanbury Brown-Twiss (HBT) method [33]. As in Ref. [34], the HBT radius can be parameterised as a function of multiplicity as:

$$R = a \langle dN_{\text{ch}}/dy \rangle^{1/3} + b \quad (1.7)$$

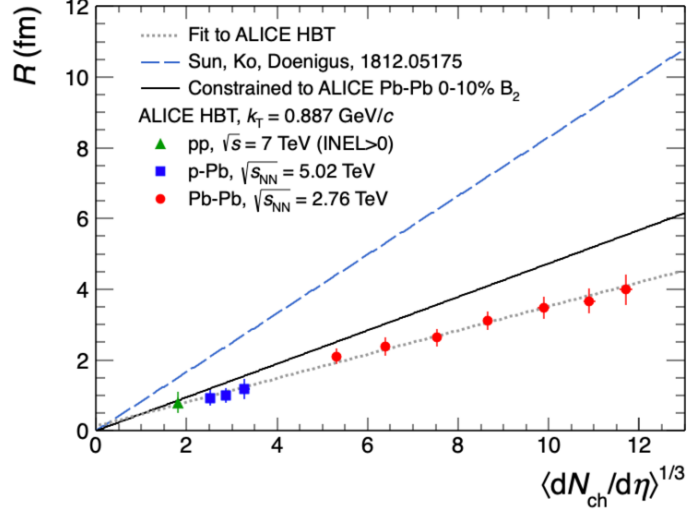


Fig. 1.7.: Comparison of different source volume parameterisations. The dotted gray line is the result of a linear fit to the ALICE HBT data [35–37], according to Eq. 1.7. The solid black line corresponds to a different parameterisation, obtained by fixing the parameters of Eq. 1.7 in order to reproduce B_2 in the most central Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV (see Sec. 6.2). The dashed blue is the relation between the system radius and the charged particle multiplicity used in Ref. [38], which does not rely on the HBT data. (Figure taken from Ref. [34]).

where the parameters a and b are obtained from a fit. Some examples of parameterisation of the HBT radius as a function of multiplicity are shown in Fig. 1.7.

The energy dependence of particle multiplicity is shown in Fig. 1.8. The yields of different hadron species are measured at mid rapidity ($|y| < 0.5$) in central HIC (Au–Au and Pb–Pb). There is a dependence of the multiplicity on the collision energy. At low energy ($\sqrt{s_{NN}} < 5$ GeV), the data measured by AGS show that the final state is dominated by the incoming nucleons (nuclear stopping regime). The difference in the yields between a particle and its respective anti-particle reflects their difference in isospin. At high energies ($\sqrt{s_{NN}} > 100$ GeV), the final state is dominated by the particle produced in the collision (nuclear transparency regime) and the same production of particles and anti-particles is observed. The energy dependence disappears at the LHC energies ($\sqrt{s_{NN}} \gg 100$ GeV), as it will be shown in the next sections.

1.8 Momentum spectra and radial flow

Low momentum hadrons are the majority ($\approx 99\%$) of the particles produced in a collision. Momentum spectra of identified particles give a snapshot of the collision

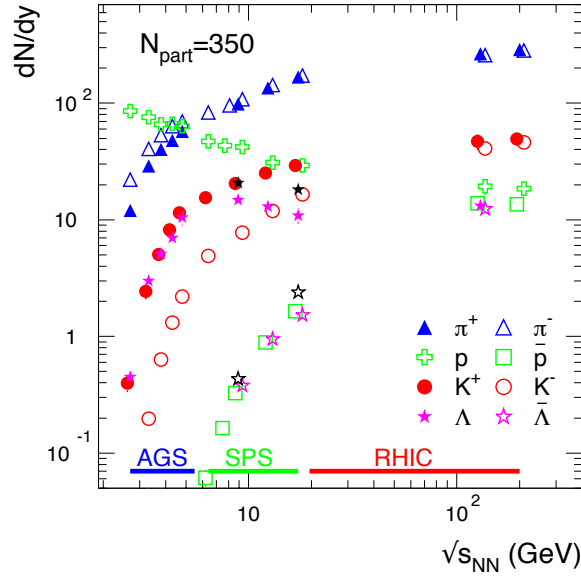


Fig. 1.8.: Energy dependence of different hadron yields at mid-rapidity ($|y| < 0.5$) as a function of collision energy [39].

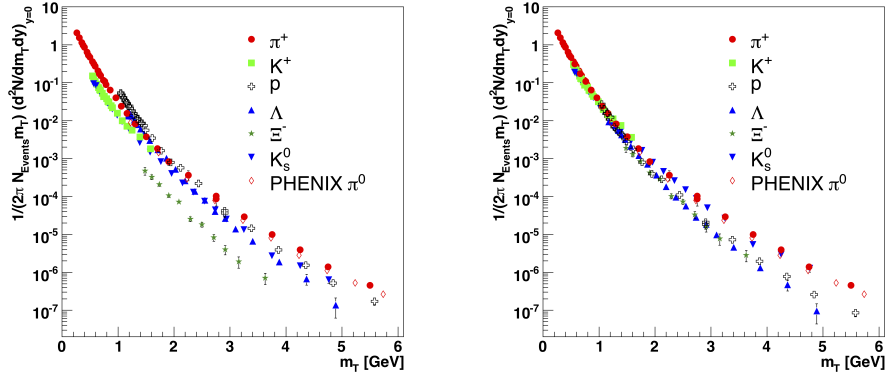
system at kinetic freeze-out. Assuming a Boltzmann-Gibbs distribution for the particle emission, the particle spectrum for the species i can be described as:

$$\frac{1}{m_T} \frac{d^2 N_i}{dm_T dy} \propto e^{-\frac{m_T}{T_{slope}}} \quad (1.8)$$

where $m_T = \sqrt{m_i^2 + p_T^2}$ is the transverse mass and T_{slope} is a fit parameter.

Keeping into account only the thermal component of the spectra and assuming a common emission temperature, the spectra of all the particle species should have the same slope. This effect, called m_T scaling, was observed in Minimum Bias (MB) pp collisions by STAR and PHENIX at $\sqrt{s} = 200$ GeV [40]. Fig 1.9a shows that the spectra for different species have the same slope for $p_T < 2$ GeV/ c and different yields, as expected. To have a better look at the slope, Fig. 1.9b shows the same spectra but properly scaled to match at $m_T = 1$ GeV/ c . The STAR collaboration measured a kinetic freeze-out temperature of 127 ± 13 MeV in Minimum Bias pp collisions at $\sqrt{s} = 200$ GeV [41]: the spectra are the Boltzmann thermal spectra and T_{slope} corresponds to the kinetic freeze-out temperature T_{kin} .

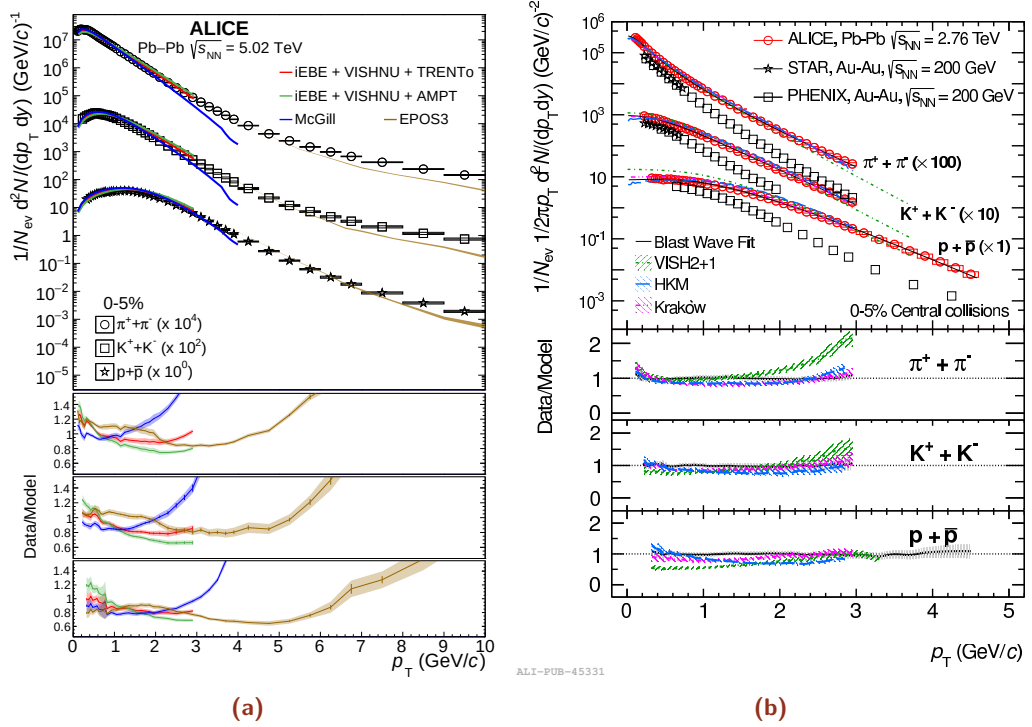
The m_T scaling is broken in HI collisions. Fig. 1.10 shows the transverse momentum spectra for pions, kaons and protons in central (0-5%) Pb-Pb collisions at $\sqrt{s_{NN}} = 5$ TeV at the LHC (Fig. 1.10a), in central (0-5%) Au-Au collisions at $\sqrt{s_{NN}} = 200$ GeV at the RHIC and in central (0-5%) Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV at the LHC (Fig. 1.10b). The slope of the spectra decreases with the mass of the particle of interest: higher mass particles are shifted towards higher values of p_T , leading



(a) Transverse mass spectra

(b) Scaled transverse mass spectra

Fig. 1.9.: Transverse mass hadron spectra at mid rapidity ($|y| < 0.5$) measured by STAR and PHENIX in pp collisions at $\sqrt{s} = 200$ GeV [40]. Error bars include systematic uncertainties.



(a)

(b)

Fig. 1.10.: Pion, kaon and proton transverse momentum spectra in central (0-5%) Pb-Pb collisions at $\sqrt{s_{NN}} = 5$ TeV [42] (a), in central (0-5%) Au-Au collisions at $\sqrt{s_{NN}} = 200$ GeV at the RHIC and in central (0-5%) Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV at the LHC [43].

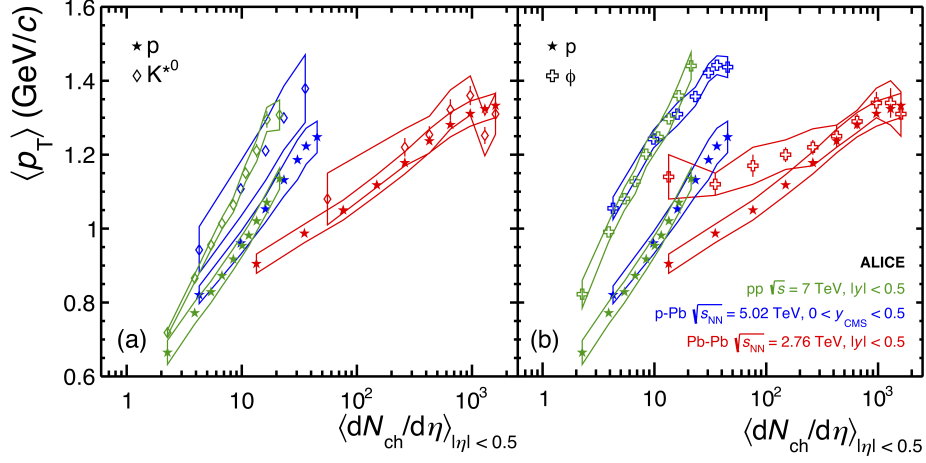


Fig. 1.11.: Mean transverse momentum $\langle p_T \rangle$ as a function of mean charged particle multiplicity $\langle dN_{ch}/d\eta \rangle$ in different collision systems and at different energies for protons and K^{*0} (a) and for protons and ϕ (b) [49].

towards a smaller value of T_{slope} . This behaviour in HI collisions can be interpreted as the superposition of a collective motion of all the particles, with mean velocity $\langle v_{\perp} \rangle$, to the thermal agitation in the transverse plane. As a result, T_{slope} can be written as:

$$T_{slope} = T_{kin} + \frac{1}{2}m\langle v_{\perp} \rangle^2 \quad (1.9)$$

The additional term reflects the mean kinetic energy acquired by the particles as an effect of a hydrodynamical collective expansion along the transverse plane, called *radial flow*. Fig. 1.10b shows also another interesting aspect: the p_T spectra measured at lower collision energy [44, 45] show a softer profile than those measured at the LHC energies [43]. An interpretation is that the radial flow is stronger at high collision energy because of the stronger pressure gradients that drive the hydrodynamical expansion at the LHC.

The radial flow interpretation is well captured by the comparison with the hydrodynamical calculations in Fig. 1.10. A proper description of the transverse momentum spectra over the full p_T range should take into account a full description of the hadronic phase following the hydrodynamical expansion (as implemented in the HKM model [46]) or some corrections related to the bulk viscosity at the chemical freeze-out (as implemented in the Krakow model [47]).

The p_T spectra of pions, kaons and protons are usually fitted simultaneously with a simplified hydrodynamical model, corresponding to a blast-wave description [48]. From these fits to the spectra one can conclude that at the LHC energies both the temperature of kinetic freeze-out T_{kin} and the radial flow velocity $\langle \beta_{\perp} \rangle$ are higher than at RHIC energies [43].

Figure 1.11 shows the mean transverse momentum $\langle p_T \rangle$ as a function of mean charged particle multiplicity $\langle dN_{\text{ch}}/d\eta \rangle$ in pp collisions at $\sqrt{s} = 7$ TeV [49], in p–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV [50] and in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [48] for protons, K^{0*} and ϕ mesons. Considering Pb–Pb collisions, it is possible to see a hardening of the spectra with increasing event multiplicity, hence with the centrality of the collision. This trend can be explained as a consequence of the hydrodynamical radial expansion of the produced medium, an effect of the increasing pressure gradients with then size of the collision system. The radial expansion is studied with the Boltzmann-Gibbs blast-wave model [51] that assumes a locally thermalised medium that expands with a common velocity field and undergoes an instantaneous freeze-out.

A similar evolution of $\langle p_T \rangle$ with multiplicity is seen in pp and p–Pb collisions. In particular, the pp data are in agreement with the p–Pb ones at similar multiplicity. It suggests that there is a mechanism that depends on the system size, regardless of the collision system. However, Pb–Pb data are consistently lower than the p–Pb ones at the same multiplicity. An explanation could be the energy loss of partons in a deconfined medium, which is produced in Pb–Pb collisions but not in smaller systems [52]. Another explanation could be that the pressure gradients are larger in small systems [53].

Finally, Fig. 1.11 shows a mass ordering: the heavier the particle, the higher the $\langle p_T \rangle$. This is another proof of radial flow: since all particles have the same mean velocity $\langle v_{\perp} \rangle$, species with higher mass have consequently a higher $\langle p_T \rangle$.

1.9 Hadrochemistry

Hadrochemistry is the study of particle species abundances. Such a study is extremely interesting for understanding the mechanisms that drive hadron production. The key observable is the p_T -integrated yield dN/dy , measured at mid rapidity ($|y| < 0.5$).

Figure 1.12 shows the ratio between the yields of different hadrons and pions, which are the most abundant species. The particle-to-pion ratio is shown as a function of multiplicity to make a direct comparison of particle yields across different collision systems. The most evident aspect of this measurement is that particle ratios evolve smoothly as a function of multiplicity, regardless of the collision system that we are considering. This will be shown to be true also for deuteron production in Sec. 6.3.

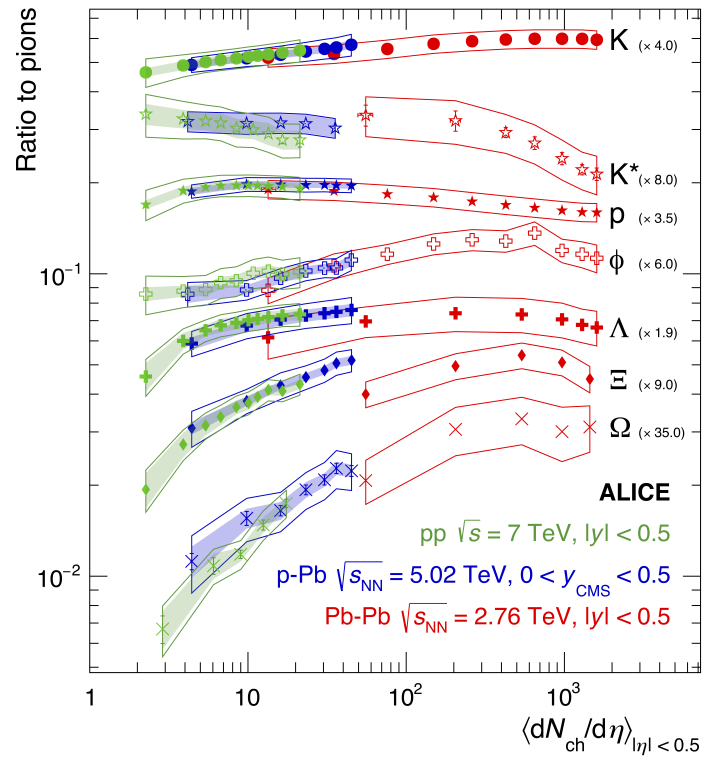


Fig. 1.12.: Integrated particle-to-pion ratios as a function of $\langle dN_{ch}/d\eta \rangle$ for pp, p-Pb and Pb-Pb collisions [49].

Concerning Fig. 1.12, the p/π ratio is observed to be constant within the uncertainties, except for the lowest multiplicity. On the contrary, Λ/π , Ξ/π and Ω/π increase with the multiplicity and the increase is more evident for hyperons with higher strangeness content. Since p/π is flat, a dependence on the mass can be excluded. However, ϕ/π increase is similar to Λ/π even if the number of strange (anti-)quarks is the double. A possible solution could be to consider also other quantum numbers, such as the baryonic number, as it will be discussed in section 2.1.2. Finally, K^{*0}/π exhibits a hint of decrease with multiplicity, more pronounced at high multiplicity. The typical interpretation in nucleus-nucleus collision is that K^{*0} decay daughters are re-scattered during the hadronic phase of the system evolution [54].

In conclusion, the relative particle abundance can be described as a universal function of charged particle multiplicity, hence of the system size.

Nuclear matter production

Light (anti-)nuclei are abundantly produced at the LHC, but their production mechanism is still not understood. Nucleons are kept together inside (anti-)nuclei by a binding energy of the order of 1 MeV per nucleon [55]. This value is extremely low if compared with the chemical freeze-out temperature of a Pb–Pb collision ($T_{ch} \sim 150$ MeV). Hence, it is surprising to see how fragile objects like light (anti-)nuclei can be produced and survive in such extreme conditions, like ‘snowballs in hell’ [56]. The production of light (anti-)nuclei has been theoretically described using two classes of phenomenological models: the Statistical Hadronisation Models (SHMs) and the coalescence models.

2.1 Statistical Hadronisation Model

The Statistical Hadronisation Model (SHM) is used to describe the abundances of hadrons produced in particle collisions. The key concept of the SHM is that particles originated from an excited region evenly occupy all the available states in phase space [57]. The idea to apply statistical concepts to the particle production was introduced by Enrico Fermi in 1950 [58] and further developed by Rolf Hagedorn [59].

In the SHM frame, also referred as Thermal Model, the final state is composed by all the particle states compatible with conservation laws imposed by the Standard Model [60] and relative particle abundances are determined from the complete filling of the phase space. Depending on the system size, two different approaches are used. On the one hand, for systems characterised by a large volume, such as Pb–Pb collisions, a grand canonical approach is used. On the other hand in small systems a canonical approach is necessary. The two approaches will be described in the following sections.

2.1.1 Grand Canonical Statistical Hadronisation Model

As discussed in Ref. [61], the grand canonical ensemble is appropriate to describe HI collisions because the reaction volume is sufficiently large. In this approach, the system under study is in contact with an energy and charge reservoir. Since the

region of interest, i.e. the mid rapidity region, makes a sufficiently small subsystem with respect to the whole interaction system, the condition is met. Therefore, system energy and charges are only conserved on average and their value is determined by the temperature and the chemical potentials associated to each charge.

The statistical properties of the system at thermodynamic equilibrium are described by the grand partition function Z :

$$Z(T, V, \mu) = \text{Tr} \left[\exp \left(-\frac{H + \sum_i \mu_i Q_i}{T} \right) \right], \quad \text{with } \mu = \sum_i \mu_i Q_i \quad (2.1)$$

where H is the Hamiltonian of the system, T is the temperature of the system, V is the system volume, Q_i are the conserved charges and μ_i their respective chemical potentials. The Hamiltonian of a Hadron Resonance Gas (HRG) is used because it can describe the interaction of a strongly interacting medium and it can reproduce the equation of state obtained with LQCD over a wide range of temperature values before the transition to a deconfined state [62]. For a strongly interacting medium, the conserved charges are the electric charge Q , the strangeness content S and the baryon number B .

The partition function Z of the system is given by the product of all partition functions Z_i of the particles included in the Hadron Resonance Gas:

$$Z(T, V, \mu) = \prod_i Z_i(T, V, \mu_i) \longrightarrow \log Z(T, V, \mu) = \sum_i \log Z_i(T, V, \mu_i) \quad (2.2)$$

The partition functions Z_i are defined as:

$$\log Z_i(T, V, \mu_i) = \frac{V g_i}{2\pi} \int_0^\infty \pm p^2 dp \log \left(1 \pm \lambda_i(T, \mu_i) e^{-\epsilon_i/T} \right) \quad (2.3)$$

where the sign $\pm$ is determined by the nature of the state, - for bosons and + for fermions. The factor g_i takes into account the degeneracy due to spin and isospin and $\epsilon_i = \sqrt{m_i^2 + p^2}$ is the energy of a particle of the species i with momentum p and mass m_i . The dependence on the chemical potential is included in the fugacity λ_i , defined as:

$$\lambda_i(T, \mu_i) = e^{(B_i \mu_B + S_i \mu_S + Q_i \mu_Q)/T} = e^{\mu/T} \quad (2.4)$$

where B_i , S_i and Q_i are respectively the baryon number, the strangeness content and the electric charge of the particle species i . As shown in Ref. [63], expanding the logarithm and performing the momentum integration one obtains:

$$\log Z_i(T, V, \mu_i) = \frac{VT g_i}{2\pi^2} \sum_{k=1}^{\infty} \frac{(\pm 1)^{k+1}}{k^2} \lambda_i^k m_i^2 K_2 \left(\frac{k m_i}{T} \right) \quad (2.5)$$

where K_2 is the modified Bessel function. From the partition function the average number $\langle N_i \rangle^{th}$ of particles of the species i is obtained through the equation:

$$\langle N_i \rangle^{th}(T, V, \mu) = T \frac{\partial}{\partial \mu_i} \log Z_i(T, V, \mu_i) = \frac{VT g_i}{2\pi^2} \sum_{k=1}^{\infty} \frac{(\pm 1)^{k+1}}{k} \lambda_i^k m_i^2 K_2 \left(\frac{k m_i}{T} \right) \quad (2.6)$$

To compare the theoretical prediction with a real measurement of the yield, it is first necessary to add the feed-down contribution of other species j (resonances) that decay in a particle of the species i :

$$\langle N_i \rangle(T, V, \mu) = \langle N_i \rangle^{th}(T, V, \mu) + \sum_j \Gamma_{j \rightarrow i} \langle N_j \rangle^{th}(T, V, \mu) \quad (2.7)$$

where $\Gamma_{j \rightarrow i}$ is the branching ratio of the species j decaying into the species i .

In this simple version, the HRG is described as a gas of not interacting hadrons and resonances. However, this condition is satisfied only in a low density system, where repulsive and attractive interactions can be neglected. The repulsive corrections, whose implementation is currently matter of study, can be modelled with eigenvolume corrections of the van der Waals type, as reported in Ref. [64].

Equation 2.7 shows that particle yields depend on five parameters: the system temperature (T), the system volume (V) and the chemical potentials (μ_B , μ_S and μ_Q). However, two chemical potentials are constrained by the initial state conditions. Indeed, μ_Q is fixed by the isospin asymmetry of the initial stage and μ_S is fixed to $\mu_S = 0$ by the null net strangeness content in the colliding nuclei. However, μ_B is not constrained by initial conditions because the net baryon content in the interaction region depends on the energy of the collision. Finally, the dependence on the volume can be removed considering ratios between the yields of different species rather than the yields themselves.

The statistical thermal model is used to predict hadron yields in a collision, going from pions to light nuclei. In principle one would not expect to see any light nuclei in HI collisions, since the chemical freeze-out temperature is much higher than their binding energy. However, the thermal model is blind to the internal structure of hadrons and the hadron yields are the result of their distribution in phase space. Using the thermal model it is possible to predict not only the energy dependence of production yields for light (anti-)nuclei, but also for multi-strange light (anti-)hypernuclei, i.e. nuclei where a neutron is replaced with a hyperon. Fig. 2.1 shows thermal model predictions for the production yields of nuclei and hypernuclei as a function of the collision energy [65]. It is possible to see that at low energy the production of nuclei (solid line) is favoured against that of antinuclei (dashed line). The reason is the dependence of the baryon chemical potential μ_B on

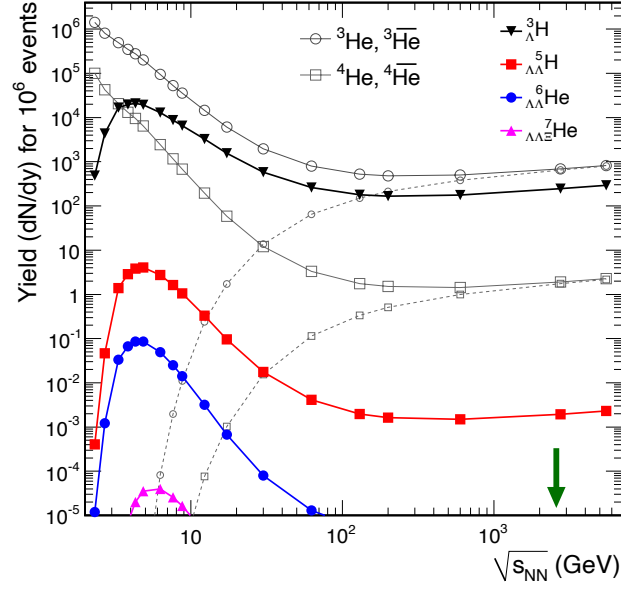


Fig. 2.1.: Thermal model predictions for the production yields of nuclei and hypernuclei as a function of the collision energy for one million Pb–Pb collisions [65]. The green arrow points to $\sqrt{s_{\text{NN}}} = 2.76$ TeV, value reached at the LHC.

the collision energy. Indeed, at low collision energies the baryon chemical potential differs from zero and the production of matter is favoured against anti-matter. If the collision energy increases the baryon chemical potential decreases, approaching zero, and the same amount of matter and anti-matter is predicted.

Figure 2.2 shows the thermal model fit to all the light-flavoured hadrons measured in central (0–10%) Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV up to hypertriton (${}^3_{\Lambda}\text{H}$). For ${}^3\text{He}$ and ${}^4\text{He}$ the predictions of the yields, computed with the parameters obtained from the fit, are reported. The fits are performed with three different implementations of the thermal model: THERMUS [66], GSI-Heidelberg [67] and SHARE [68]. These models differ in the list of resonances included in the Hamiltonian. Moreover, while hadrons are assumed to be point-like in SHARE, the hadron volume is taken into account in THERMUS and in GSI-Heidelberg. For all the models, the fit to the hadron yields provide a chemical freeze-out temperature $T_{\text{chem}} = 156$ MeV [69]. The measured yields of ${}^4\text{He}$ and ${}^4\overline{\text{He}}$ agree within the uncertainties with a temperature between 137 MeV and 177 MeV, depending on the model. This means that light nuclei, from (anti-)deuterons to (anti-) ${}^4\text{He}$, are produced at statistical equilibrium at the same temperature as the other hadrons.

2.1.2 Canonical Statistical Hadronisation Model

In systems characterised by a small reaction volume, the conditions for a grand canonical approach are not satisfied. Indeed, in small systems such as pp collisions,

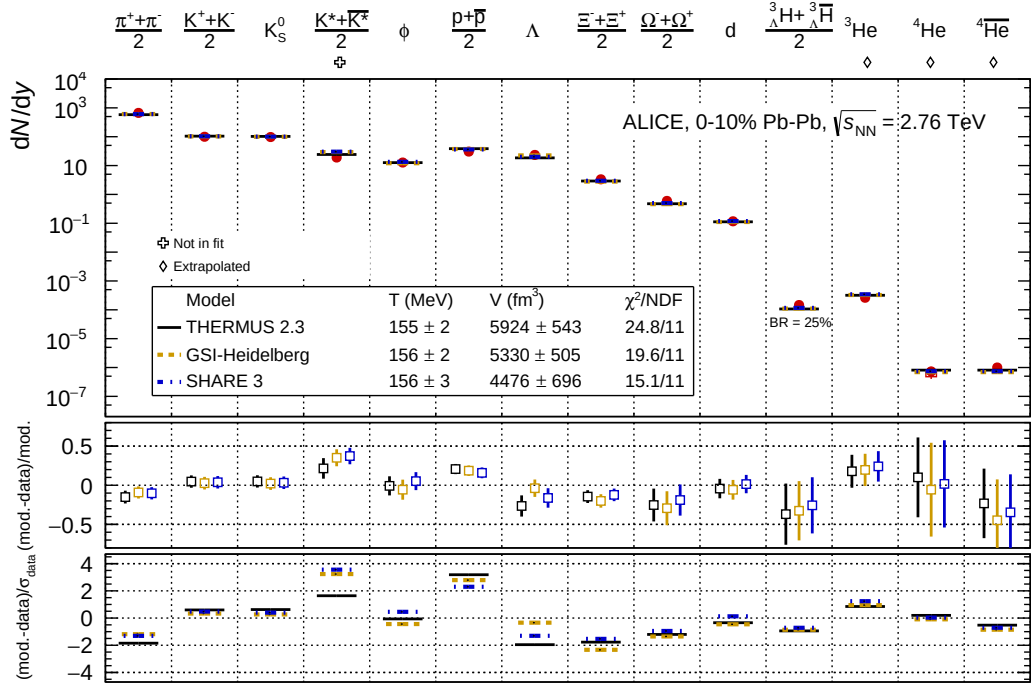


Fig. 2.2.: Thermal model fit to the light-flavoured hadron yields in central Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [69], with three different SHM implementations: THERMUS (black line), GSI-Heidelberg (yellow line) and SHARE (blue line).

in which the number of particles with a conserved charge is of the order of unity or smaller, the local conservation of charges must be constrained [70].

While in a grand-canonical ensemble all conserved charges are conserved on average, in a canonical ensemble the charges are conserved exactly and they do not fluctuate from one microscopic state to another. As a result, the yields of the particles carrying the conserved charges are suppressed with respect to a grand canonical description. This effect is called *canonical suppression*. The canonical statistical models (CSMs) have been successfully used to describe hadron yields measured in small systems, such as e^+e^- collisions [71], pp or $p\bar{p}$ collisions [72].

In this section the CSM implementation of Ref. [73] is considered. This model implements the full canonical treatment of baryon number, electric charge and strangeness content and includes light nuclei. Before this study other CSM versions were used to describe hadron abundances at the LHC energies, but they required only the local conservation of the strangeness content [74] or they did not take into account light nuclei [75]. The exact conservation of all the three conserved charges has an important effect on particle abundances. In particular, nuclei are an extremely interesting test for canonical suppression because they are characterised by a baryon number greater than the unity and hence they are affected in larger measure by canonical suppression.

Similarly to the SHM described in section 2.1.1, the system is described as an ideal hadron resonance gas in full thermal equilibrium. In the following the Boltzmann approximation will be used. The three conserved charges (Q , S and B) are fixed exactly to the values which are conserved across the correlation volume V_c . The partition function at a given temperature T and correlation volume V_c has the form:

$$Z(B, Q, S) = \int_{-\pi}^{\pi} \frac{d\phi_B}{2\pi} \int_{-\pi}^{\pi} \frac{d\phi_Q}{2\pi} \int_{-\pi}^{\pi} \frac{d\phi_S}{2\pi} e^{-i(B\phi_B + Q\phi_Q + S\phi_S)} \times \exp \left[\sum_j z_j^1 e^{i(B_j\phi_B + Q_j\phi_Q + S_j\phi_S)} \right] \quad (2.8)$$

where the sum is over all the particle species j included in the hadron resonance gas and B_j , Q_j and S_j are the charges carried by the species j . The single-particle partition function z_j^1 for the species j is defined as:

$$z_j^1 = V_c \int dm \rho_j(m) g_j \frac{m^2 T}{2\pi^2} K_2 \left(\frac{m}{T} \right) \quad (2.9)$$

where g_j is the degeneracy factor for the species j , $\rho_j(m)$ is the mass distribution of the species j , which takes into account the finite width of resonances, K_2 is the modified Bessel function.

Following the method shown in Ref. [76], the mean particle multiplicity for the species j calculated in the canonical ensemble $\langle N_j^{th} \rangle^{ce}$ can be written as:

$$\langle N_j^{th} \rangle^{ce} = \frac{Z(B - B_j, Q - Q_j, S - S_j)}{Z(B, Q, S)} \langle N_j^{th} \rangle^{gce} \quad (2.10)$$

where $\langle N_j^{th} \rangle^{gce}$ is the mean particle multiplicity for the species j calculated in the grand canonical ensemble at the same temperature T and at the same volume V_c . The factor before $\langle N_j^{th} \rangle^{gce}$ is the chemical factor, which reflects the exact conservation of the charges. As for Eq. 2.7, the final particle yield $\langle N_j \rangle^{ce}$ is obtained by adding the feed-down contributions from resonances:

$$\langle N_j \rangle^{ce} = \langle N_j^{th} \rangle^{ce} + \sum_i \Gamma_{i \rightarrow j} \langle N_i^{th} \rangle^{ce} \quad (2.11)$$

where $\Gamma_{i \rightarrow j}$ is the branching ratio of the species i decaying into the species j .

In Ref. [73] the CSM is implemented within the *Thermal-FIST* package [77]. One of the key features of this package is that the conserved charges are fixed to the initial values $Q = S = B = 0$. Indeed, at the LHC energies the net baryon content at mid rapidity is null, as shown by the anti-baryon/baryon ratios compatible with

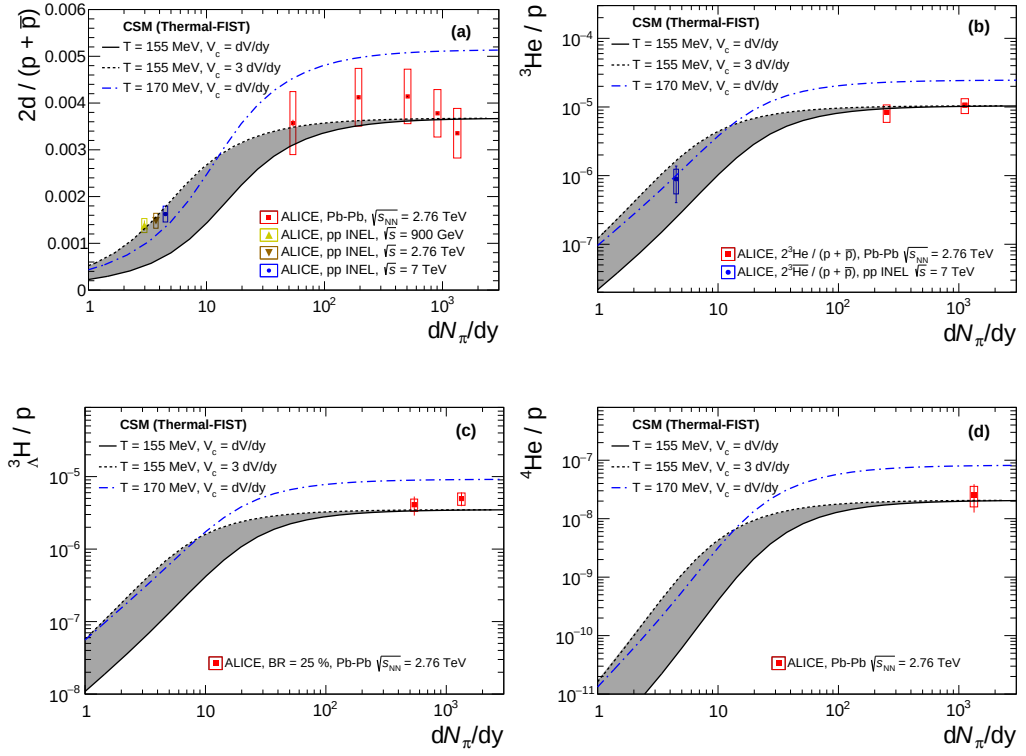


Fig. 2.3.: Predictions for d/p (a), ${}^3\text{He}/p$ (b), ${}^3_\Lambda\text{He}/p$ (c) and ${}^4\text{He}/p$ (d) as a function of charged pion multiplicity dN_π/dy obtained with *Thermal-Fist* package. Predictions are reported for $T = 155$ MeV and $V_c = dV/dy$ (black solid lines), $T = 155$ MeV and $V_c = 3 dV/dy$ (black dashed lines) and $T = 170$ MeV and $V_c = dV/dy$ (blue dash-dotted lines). Experimental points corresponds to the available ALICE measurements at the time of the publication of Ref. [73], from which this picture is taken.

unity [78]. In this implementation of CSM, all the particles in the HRG are considered point-like. Considering a finite volume (exclusion volume) for all the particles would affect the yields and for this reason a new implementation of the model that takes into account exclusion volumes is currently under study. The temperature of chemical freeze-out is fixed to $T_{ch} = 155$ MeV, in agreement with the fits to the production yields of hadrons in central Pb-Pb collisions performed by the ALICE experiment [69]. The value of the correlation volume V_c is varied to study the multiplicity dependence of the production.

Figure 2.3 shows the prediction of the ratios between the p_T -integrated yields dN/dy of different nuclei (and of hypertriton) and that of protons as a function of pion multiplicity dN_π/dy . The pion multiplicity dN_π/dy is used by the authors because it is measured in the experiments and it can be calculated in the CSM. As reported in Eq. 2.11, CSM computes $\langle N_j \rangle^{ce}$, but experiments measure dN/dy . Therefore, it is necessary to convert the former into the latter. The conversion factor is obtained from the connection between the correlation volume V_c and the volume dV/dy

corresponding to one unity of rapidity. In Ref. [79], arguments related to causal connections within the expanding system limit the correlation volume to the total system volume and suggest a value of few rapidity units. The conversion can be expressed as:

$$V_c = k \cdot dV/dy \longrightarrow dN_j/dy = \langle N_j \rangle^{ce}/k \quad (2.12)$$

Therefore a variation in the conversion parameter k has an important effect on the yields. In Fig. 2.3, CSM predictions are reported for a correlation volume $dV/dy < V_c < 3 dV/dy$ and for a temperature $T = 155$ MeV. Experimental data include the yield ratios $2d/(p + \bar{p})$, $2^3\text{He}/(p + \bar{p})$, $^3_\Lambda\text{H}/p$ and $^4\text{H}/p$ measured in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [69, 80, 81], the yield ratio $2d/(p + \bar{p})$ measured in inelastic pp collisions at $\sqrt{s} = 0.9, 2.76$ and 7 TeV [82], and the yield ratio $2^3\overline{\text{He}}/(p + \bar{p})$ measured in inelastic pp collisions at $\sqrt{s} = 7$ TeV [82]. The theoretical predictions grow monotonically with increasing dN_π/dy and eventually saturate at the corresponding grand-canonical values at high multiplicity. The few experimental points available at low multiplicity seem to confirm canonical suppression. Moreover, a better agreement is shown for $V_c \gtrsim 3 dV/dy$. A prediction for $T = 155$ MeV and $V_c = dV/dy$ is also reported, but it overestimates experimental data.

In Sec. 6.3, the CSM predictions for $2d/(p + \bar{p})$ as a function of charged particle multiplicity $dN_{\text{ch}}/d\eta$ will be shown for the latest published ALICE results and for this work.

2.2 Coalescence model

Hadron coalescence models assume that nucleons, i.e. protons and neutrons, that are close to each other in phase space after chemical freeze-out can bind together and form a nucleus via coalescence. As well as SHMs, coalescence models do not provide a detailed description of the interaction that leads to nucleus formation.

The coalescence model was first developed by Butler and Pearson for proton-nucleus collisions [83]. In their study, they suggested that the deuteron production mechanism was the binding of cascade nucleons in the presence of the target nuclear optical potential.

The key parameter of coalescence model is the *coalescence parameter* B_A , defined as:

$$B_A(p_T^p) = E_A \frac{d^3 N_A}{dp_A^3} \bigg/ \left(E_p \frac{d^3 N_p}{dp_p^3} \right)^A = \frac{1}{2\pi p_T^A} \frac{d^2 N_A}{dy dp_T^A} \bigg/ \left(\frac{1}{2\pi p_T^p} \frac{d^2 N_p}{dy dp_T^p} \right)^A. \quad (2.13)$$

The numerators of the second and of the third terms of Eq. 2.13 are two equivalent forms of the invariant production spectra of nuclei with mass number A , while the denominators are the invariant production spectra for primary protons, i.e. produced in a collision, elevated to the mass number A . In principle one should consider proton and neutron invariant spectra separately. However, neutrons and protons belong to the same isospin multiplet and at the LHC energies the isospin chemical potential is expected to be vanishing. Therefore, the production spectra of protons and neutrons are the same. In addition, in several experiments, such as ALICE, neutrons cannot be detected because of their null electric charge and to the absence of a hadronic calorimeter. Moreover, in Eq. 2.13 the invariant spectrum of protons is evaluated at $p_T^p = p_T^A/A$, because the momentum of the nucleus is the sum of the momenta of A nucleons. For what concerns the work presented in this thesis, which is related to the production of (anti-)deuterons, the mass number is $A = 2$ and the computed coalescence parameter is B_2 .

One of the first implementations of the coalescence model is the *simple* coalescence model, also referred as *naive* coalescence model [84]. The reason for such a nomenclature is that in this implementation only momentum-space is considered and not space-time: for two nucleons to coalesce it is sufficient that they have similar momenta, i.e. they are close in momentum space, regardless of their space distance. In small systems, such as pp collisions, this assumption is plausible because the system size is comparable with the size of the nucleus, thus nucleons are produced close to each other. On the contrary, in large systems, such as central Pb–Pb collisions, the system size is larger than the nucleus size and two nucleons can be produced very far from each other and hence they cannot coalesce into a nucleus even though their momenta are similar. The system size can be determined from the charged-particle multiplicity (see Sec. 1.7). Neglecting nucleon isospin, the coalescence parameter can be expressed as:

$$B_A = \left(\frac{4\pi}{3} p_0^3 \right)^{A-1} \frac{m_A}{m_p^A} \quad (2.14)$$

where A is the nucleus mass number, m_A is the nucleus mass, m_p the proton mass and p_0 is the maximum distance in momentum-space at which coalescence can happen. According to the naive coalescence model, B_A should not depend neither on transverse momentum nor on the system volume, hence on the charged particle multiplicity. However, as shown in Fig. 2.4, even if B_A is flat as a function of

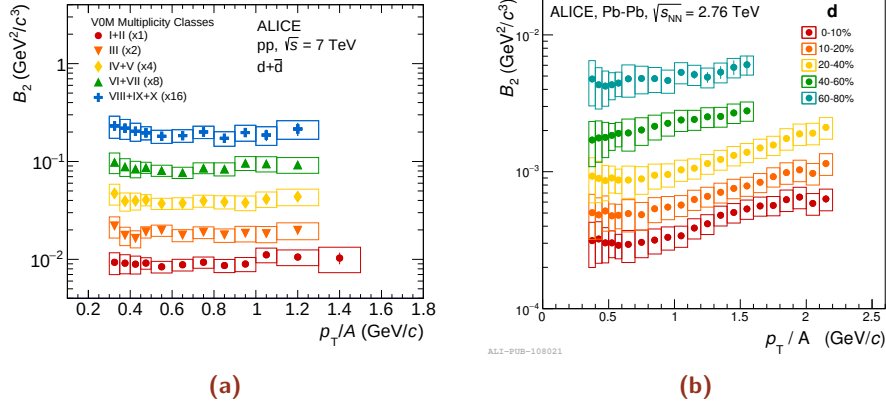


Fig. 2.4.: Measured B_2 in pp collisions at $\sqrt{s} = 7 \text{ TeV}$ [85] (a) and in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76 \text{ TeV}$ [80] (b).

transverse momentum in pp collisions [85], in large systems such as Pb-Pb collisions B_A increases with transverse momentum [80]. Similarly, as it will be shown in Sec. 6.2, B_A decreases with charged particle multiplicity. Indeed, the coalescence probability decreases for nucleons with similar momentum that are distant in space. For this reason, a more advanced implementation of the coalescence model that takes into account space-time must be considered to correctly reproduce the trend observed in the data.

In Ref. [34], a more advanced implementation of the coalescence model, which takes into account the system size, is described. The key concept is the overlap between the phase-space distributions of the nucleons and the phase space of the nucleus. Due to the quantum-mechanical nature of coalescence, the phase space of the nucleus is replaced by its Wigner function, obtained via the Wigner transform of its wave function.

From here on the attention is focused on the deuteron. The deuteron wave-function is approximated by the ground-state of an isotropic harmonic oscillator with characteristic-size parameter r_d :

$$\varphi_d(\vec{r}) = (2\pi r_d^2)^{-3/4} \exp\left(-\frac{r^2}{2r_d^2}\right) \quad (2.15)$$

The characteristic-size parameter is defined as $r_d = \sqrt{8/3} \lambda_d$, where λ_d is the rms of the charge wave-function of the deuteron. The harmonic-oscillator approximation has the advantage that the resulting Wigner function is gaussian and allows for fully analytical solutions.

As shown in Ref. [86], the dependence of B_2 on the source volume is expressed by the equation:

$$B_2 = \frac{3\pi^{3/2}\langle C_d \rangle}{2m_T R_\perp^2(m_T) R_\parallel(m_T)} \quad (2.16)$$

where $R_\perp$ and $R_\parallel$ are the homogeneity radii of the source, which can be measured with HBT interferometry [33]. The factor $\langle C_d \rangle$ is the quantum-canonical correction factor, which accounts for the quantum-canonical nature of the coalescence process, and can be approximated as:

$$\langle C_d \rangle \approx \frac{1}{\left[1 + \left(\frac{r_d}{2R_\perp(m_T)}\right)^2\right] \sqrt{1 + \left(\frac{r_d}{2R_\parallel(m_T)}\right)^2}} \quad (2.17)$$

Assuming $R_\perp \approx R_\parallel \approx R$, Eq 2.16 and Eq. 2.17 simplify to:

$$B_2 = \frac{3\pi^{3/2}\langle C_d \rangle}{2m_T R^3(m_T)} \quad (2.18)$$

and

$$\langle C_d \rangle \approx \left[1 + \left(\frac{r_d}{2R(m_T)}\right)^2\right]^{-3/2} \quad (2.19)$$

The evolution of B_2 and $\langle C_d \rangle$ with the source volume is shown in Fig. 2.5, with different values of the characteristic-size parameter r_d . For a point-like deuteron ($r_d = 0$), the quantum-mechanical correction factor is identically one, as expected. For finite values of r_d , $\langle C_d \rangle$ is identically null for a point-like source, increases with increasing source radius R and eventually saturates to unity. The larger is r_d , the slower is the raise. The presence of $\langle C_d \rangle$ in the definition of B_2 is crucial. For low R values, B_2 is dominated by the correction factor, with a finite value for $R \rightarrow 0$. For $R \gg r_d$, the coalescence probability is dominated by the classical phase-space separation and decreases for large distances in configuration space. For the comparison with experimental data, $r_d = 3.2$ fm is used. This value corresponds to the measured rms radius of the charge wave-function of deuteron, $\lambda_d = 2.1413 \pm 0.0025$ fm [87].

Also in Ref. [38] the coalescence model is implemented using the Wigner formalism. Here the deuteron wave-function is approximated by an isotropic harmonic oscillator with characteristic-size parameter $r_d = \sqrt{8/3} \lambda_d$, with $\lambda_d = 1.96$ fm [88]. The phase-space distributions of nucleons are taken from a thermalised expanding spherical fireball of kinetic freeze-out temperature T_k and radius R :

$$f_{p,n}(\vec{x}, \vec{p}) = \frac{N_{p,n}}{(2\pi)^3 (mT_k R^2)^{3/2}} e^{-\frac{p^2}{2mT_k}} e^{-\frac{x^2}{2R^2}} \quad (2.20)$$

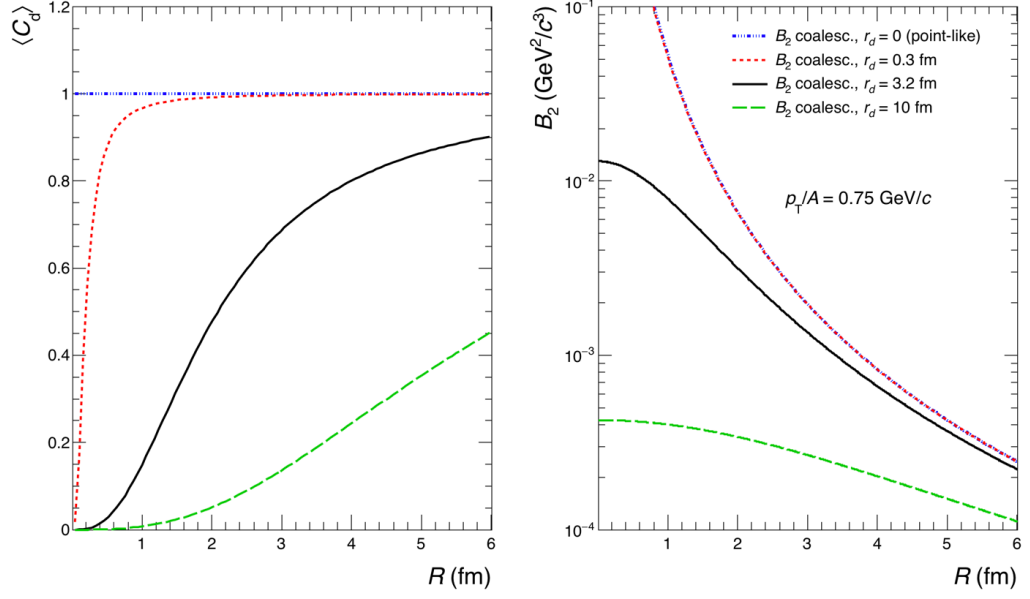


Fig. 2.5.: Quantum-mechanical correction factor $\langle C_d \rangle$ (left) and coalescence parameter B_2 (right) as a function of the radius R of the source. Both the quantities are computed assuming different size parameter r_d for the size of the deuteron. The figure is taken from Ref. [34]

where m is the mass of the nucleon and $N_{p,n}$ is the number of protons or neutrons. From the full calculation one obtains:

$$N_d = \frac{3N_p N_n}{4(mT_k R^2)^{3/2}} \frac{1}{\left(1 + \frac{1}{mT_k r_d^2}\right)^{3/2}} \frac{1}{\left(1 + \frac{r_d^2}{4R^2}\right)^{3/2}} \quad (2.21)$$

Since the kinetic freeze-out temperature T_k for nucleons is of the order of 100 MeV, $mT_k \gg 1/r_d^2$ and hence the yield ratio d/p can be expressed as:

$$\frac{N_d}{N_p} \approx \frac{3N_n}{4(mT_k R^2)^{3/2}} \left[1 + \left(\frac{1.6 \text{ fm}}{R}\right)^2\right]^{-3/2} \quad (2.22)$$

The last factor on the right-hand side of Eq. 2.22 expresses the suppression of deuteron production, due to its finite size with respect to the nucleon emission source. Its value is significantly lower than unity for a radius $R < 1.6$ fm and approaches unity when the radius R is larger than the deuteron size. The first factor on the right-hand side of Eq. 2.22, instead, corresponds to the d/p limit for a source with radius R much larger than the deuteron size.

The evolution d/p as a function of multiplicity, hence of the system size, is reported in panel (a) of Fig. 2.6. It is possible to see that for low multiplicities, where the system size is smaller or comparable with the size of the deuteron, the second factor in Eq. 2.22 is dominant. For large values of the system size the second factor approaches unity and d/p ratio is determined by the first term. For the comparison

with data, the neutron number N_n can be replaced with the proton number N_p , due to the vanishing isospin chemical potential at the LHC. Finally, T_k is extracted from the measured charged particle spectra.

In Chapter 6, the measurements of B_2 and d/p performed for the present thesis project will be compared with the predictions of the theoretical models shown above.

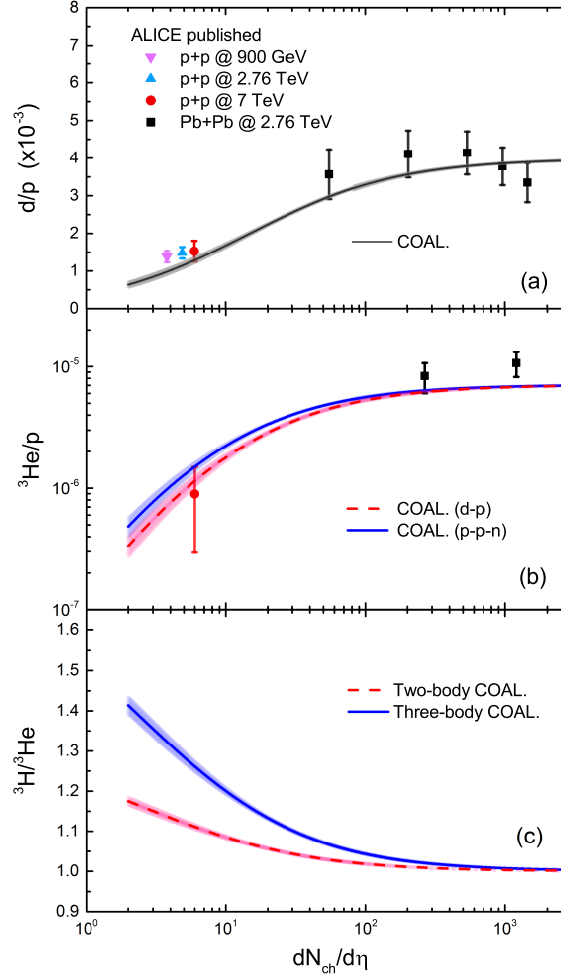


Fig. 2.6.: Yield ratios d/p , $^3\text{He}/p$ and $^3\text{H}/^3\text{He}$ as a function of charged particle multiplicity. Experimental data are measured in pp collisions at $\sqrt{s} = 0.9, 2.76$ and 7 TeV [80] and in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [89]. Lines are the predictions of the coalescence model, with theoretical uncertainties on the emission source radius given by the shaded bands. For $^3\text{He}/p$ and $^3\text{H}/^3\text{He}$ the magenta and the blue lines correspond to two-body and three-body coalescence, respectively. The figure is taken from Ref. [38]

2.3 Inputs for dark matter search and QCD bound state understanding

The knowledge of the production mechanisms of light (anti-)nuclei would provide inputs for the background determination in dark matter searches in space and for the understanding of the formation of QCD bound states in high energy hadron physics.

Some dark matter models predict that dark matter annihilation can produce low energy anti-deuterons [90, 91]. The possible presence of anti-deuterons in our galaxy could hence be explained as the result of dark matter annihilation [92] or as the product of the interaction of high-energy cosmic rays with the interstellar medium [93]. The contributions from cosmic rays interactions at low energy (1-2 GeV/ c per nucleon) should be small [90, 91, 94, 95] and the sensitivity to dark matter annihilation should be enhanced. Even if in the last years the search of such a signal has been intensified with the experiments AMS-02 [96] and GAPS [97], no clear evidence of anti-nuclei production was found [98, 99]. Dedicated experiment for the search of anti-deuteron and anti-helium from dark matter annihilation are ongoing [100, 101]. A realistic model for the production of anti-deuterons would provide an input for the determination of background in the search for dark matter.

In the last decades, exotic multi-quark states have been observed in high-energy physics experiments [102]. The nature of such states is still not clearly understood. One example of these exotic states is given by the $X(3872)$, which was observed for the first time by the Belle experiment in 2003 [103]. Its existence was then confirmed by other experiments [104–106] but it is still not clear if it is a hadronic molecule, a compact tetraquark, a charmonium state, or their mixture. Indeed, the proximity of the $X(3872)$ mass to the sum of D^0 and $\bar{D}^{*0}$ masses led to a possible interpretation of $X(3872)$ as a loosely bound $D^0\bar{D}^{*0}$ state [107, 108]. A better understanding of the production mechanisms of (anti-)deuterons, which are genuine hadronic molecules, could shed light on the formation of exotic QCD bound states and eventually confirm or exclude that they are hadron molecules.

A Large Ion Collider Experiment (ALICE)

ALICE is one of the four major experiments installed at the Large Hadron Collider (LHC) of the CERN laboratory (Geneva, Switzerland). The LHC is at present the most powerful particle collider in the world, reaching colliding energies in the TeV regime. In this context the ALICE experiment plays a crucial role in the study of the properties of strongly-interacting matter and soft QCD observables in hadronic collisions.

3.1 The Large Hadron Collider

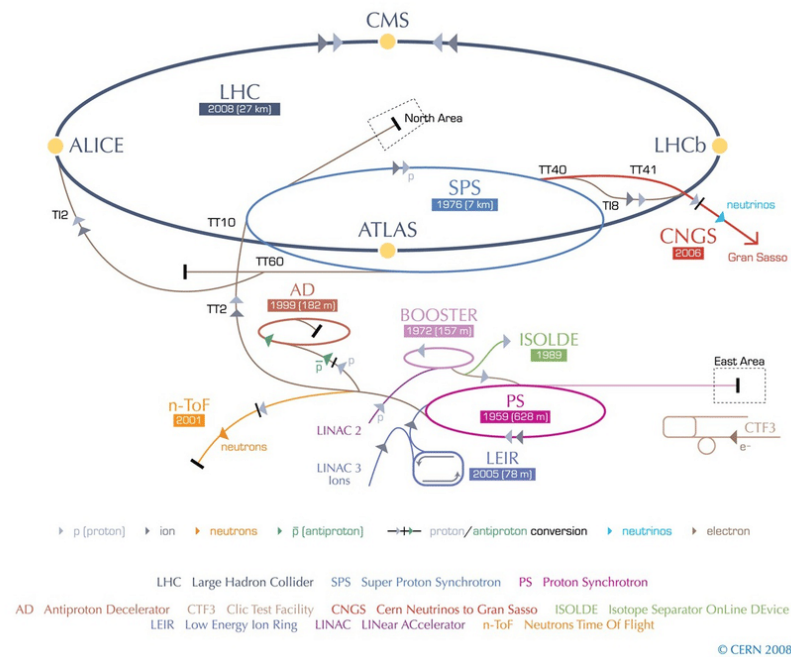


Fig. 3.1.: Scheme of the accelerator complex at CERN [109].

The LHC is the last step of the complex accelerator system at CERN (Fig. 3.1) [109]. Each element of the chain accelerates the particles to increasingly higher energies, until they reach the LHC. At CERN both protons and Pb nuclei are accelerated, but their path from source to collision is different.

Protons are extracted from an ionised-hydrogen tank and injected in the linear accelerator LINAC2, where they reach an energy of 50 MeV. Then, they are accelerated to 1.4 GeV in the Proton Synchrotron Booster (PSB) and injected in the Proton Synchrotron (PS), where they reach an energy of 25 GeV. At this stage the protons are structured in bunches. Protons are injected from the PS into the Super Proton Synchrotron (SPS), where they are accelerated to 450 GeV before their final injection into the LHC.

The initial stages of the acceleration of Pb nuclei are different. ^{208}Pb ions are extracted from a purified sample of lead, brought to high temperatures ($\sim 500^\circ\text{C}$) with microwaves. The gas is partially ionised by a stream of electrons and only Pb^{29+} are selected. Then the ions travel through the LINAC3, where they are accelerated to 4.5 MeV per nucleon and are injected into the Low Energy Ion Ring (LEIR). Here the beam is split into four bunches and brought to 72 MeV per nucleon. From the LEIR, the ion beam is injected into the PS and follows the same steps described for the protons above.

In the LHC two beams circulate in opposite directions and collide in the four interaction points corresponding to the major LHC experiments: ATLAS, ALICE, CMS and LHCb.

The LHC Run1 physics program started in 2009 delivering pp collisions up to $\sqrt{s} = 7$ TeV, p-Pb at $\sqrt{s_{\text{NN}}} = 5.02$ TeV and Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV. The Run1 ended in 2013 and the Long Shutdown 1 (LS1) started. During the LS1, the LHC has been upgraded in order to increase the collision energy. The LHC Run2 physics program started in 2015 and delivered Pb-Pb collisions at $\sqrt{s} = 5.02$ TeV, p-Pb collisions at $\sqrt{s} = 8$ TeV and pp collisions at $\sqrt{s} = 13$ TeV. In December 2018 the LHC Long Shutdown 2 started, allowing the detectors to upgrade in view of LHC Run3 physics program, that will start in 2021.

3.2 ALICE: detector overview

ALICE [110] is a general purpose heavy-ion experiment at the LHC. Its main goal is the study of the QGP and for this purpose the detector must be able to reconstruct the trajectories (tracks) of all the particles produced in a collision and to identify them. ALICE was designed when RHIC data were not available and the estimated multiplicity for heavy ion collisions at the LHC energy was in the range of 2000 - 8000 charged particles per rapidity unit [111, 112]. Consequently, relatively slow detectors with high granularity and low material budget, such as the Time Projection Chamber (TPC) and the Silicon Drift Detector (SDD), were used. The use

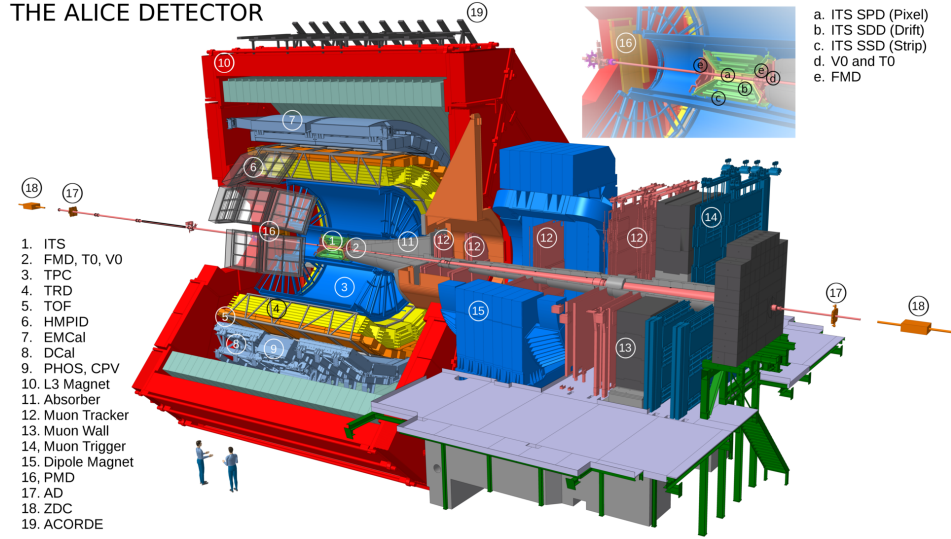


Fig. 3.2.: The ALICE experimental apparatus in LHC Run2.

of these detectors limits the maximum readout rate to 1 kHz, regardless the collision system. When RHIC measurements became available in 2001, they showed that the multiplicity range at LHC should have been 1500 - 4000 charged particles for rapidity unit and the experiment was optimised for $dN_{ch}/dy \sim 4000$.

The ALICE layout is shown in Fig. 3.2 and the details of its nineteen detectors are listed in Tab. 3.1 . The geometrical details are expressed in the coordinate system described in Sec.1.6.

The ALICE apparatus is composed of three main parts: the central barrel, which covers the mid-rapidity region ($|\eta| < 0.9$), the muon arm, which covers the forward pseudorapidity range $-4 \leq \eta \leq -2.5$, and the forward detectors.

All the detectors in the central barrel are immersed in a magnetic field of 0.5 T, generated by the ALICE solenoid magnet, which was previously used in the L3 experiment at the Large Electron-Positron (LEP) collider. The mild intensity of the magnetic field has been chosen to allow the reconstruction of the particle trajectories (tracking) of low-momentum particle in the TPC, down to $p_T \sim 80$ MeV [113]. The detectors of the central barrel, starting from the beam line and going outwards, are: the Inner Tracking System (ITS) [114], the TPC [115, 116], the Transition Radiation Detector (TRD) [117], the Time Of Flight (TOF) [118], the High Momentum Particle IDentification detector (HMPID) [119], the Photon Spectrometer (PHOS) [120], the ElectroMagnetic Calorimeter (EMCal) [121], the Di-jet calorimeter (DCal) [122] and the Charged Particle Veto detector (CPV) [123]. All these detectors are used for the tracking and the PID of particles produced at mid rapidity. In the same rapidity

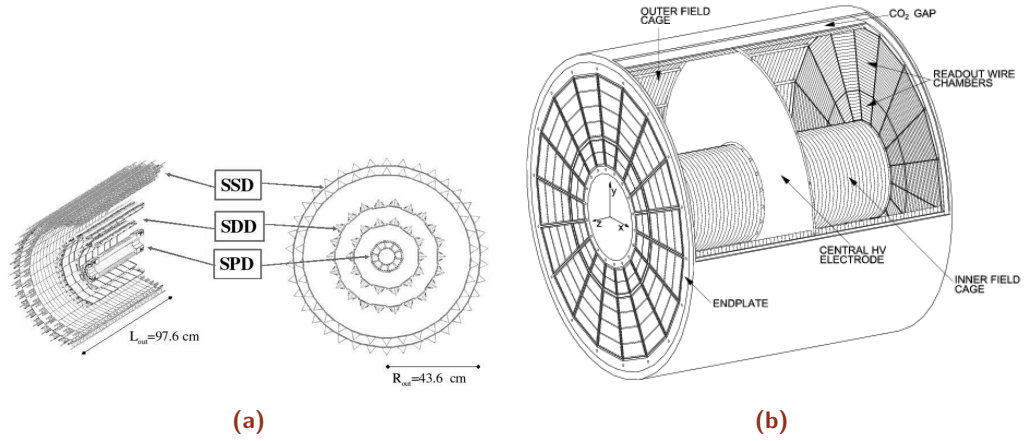


Fig. 3.3.: Layouts of the ITS (a) and of the TPC (b).

region of the central barrel but outside the magnet, one can find A Cosmic Ray Detector (ACORDE) [124], used for the study of cosmic rays.

In the muon arm one can find the Muon Spectrometer, the Muon Wall and the Muon Trigger[125]. These detectors are used to study the production of quarkonia at forward rapidity through their decay into muons.

Finally the forward detectors, which are located at a small angle with respect to the beam line, include the Zero Degree Calorimeter (ZDC) [126], the Photon Multiplicity Detector (PMD) [127], the Forward Multiplicity Detector (FMD) [128], the VZERO detector (V0) [129] and the TZERO (T0) [130]. While the ZDC is used for the measurement of the centrality of Pb–Pb collision, the PMD and the FMD are used to study the multiplicity of photons and charged particles, respectively. Finally, the V0 is a trigger detector and is used to measure the charged particle multiplicity, while the T0 provides the collision time.

The main detectors used for this analysis are described in detail in the following.

3.2.1 Inner Tracking System

The ITS is the closest detector to the collision point, located just outside the 800 μm Beryllium beam pipe. The ITS layout is shown in Fig. 3.3b. It is composed of six layers of silicon detectors based on different technologies: the two innermost layers are silicon pixel detectors (SPD), the two central layers are silicon drift detectors (SDD) and the two outermost layers are double-sided silicon strip detectors (SSD). The SPD covers a larger pseudorapidity region, $|\eta| < 1.2$ for particles generated at $z = 0$, to provide, together with the FMD, continuous coverage for the measurement of charged particle multiplicity.

Tab. 3.1.: Geometrical details about the configuration of the ALICE detectors. This table has been adapted from the extensive description of the ALICE apparatus in Ref. [131]. When it is not specified, the azimuthal coverage for the detector is 2π . The position of the detector corresponds to the radial distance from the beam axis for the barrel detectors and to the distance along z for the others. When more than one position values are specified the detector is subdivided in two or several parts and those values are the minimum and maximum distances from the interaction point.

Detector	Acceptance (η, ϕ)	Position (m)	Surface (m ²)	Channels
ITS layer 1,2 (SPD)	$\pm 2, \pm 1.4$	0.039, 0.076	0.21	9.8 M
ITS layer 3,4 (SDD)	$\pm 0.9, \pm 0.9$	0.150, 0.239	1.31	133 K
ITS layer 5,6 (SSD)	$\pm 0.97, \pm 0.97$	0.380, 0.430	5.00	2.6 M
TPC	± 0.9 at $r = 2.8$ m ± 1.5 at $r = 1.4$ m	0.848, 2.466	Readout 32.5 m ² Volume 90 m ³	557568
TRD	± 0.84	2.90, 3.68	716	1.2 M
TOF	± 0.9	3.78	141	157248
HMPID	$\pm 0.6, 1.2 \leq \phi \leq 58.8$	5	11	161280
PHOS	$\pm 0.12, 220 \leq \phi \leq 320$	4.6	8.6	17920
EMCAL	$\pm 0.7, 80^\circ \leq \phi \leq 187^\circ$	4.36	44	12672
ACORDE	$\pm 1.3, -60^\circ \leq \phi \leq 60^\circ$	8.5	43	120
Muon Tracking	$-2.5 \leq \eta \leq -4.0$	-14.22, -5.36	95	1.08 M
Muon Trigger	$-2.5 \leq \eta \leq -4.0$	-17.12, -16.12	138	21000
ZDC: ZN	$ \eta \leq 8.8$	± 116	2×0.0049	10
ZDC: ZP	$6.5 \leq \eta \leq 7.5$ $-9.7^\circ \leq \phi \leq 9.7^\circ$	± 116	2×0.027	10
ZDC: ZEM	$4.8 \leq \eta \leq 5.7,$ $-16^\circ \leq \phi \leq 16^\circ$ and $164^\circ \leq \phi \leq 196^\circ$	7.25	2×0.0049	2
PMD	$2.3 \leq \eta \leq 3.7$	3.64	2.59	221184
FMD disc 1	$3.62 \leq \eta \leq 5.03$	3.2		
FMD disc 2	$1.7 \leq \eta \leq 3.68$	0.752, 0.834	0.266	51200
FMD disc 3	$-3.4 \leq \eta \leq -1.7$	-0.752, -0.628		
VOA	$2.8 \leq \eta \leq 5.1$	3.4	0.548	32
VOC	$-3.7 \leq \eta \leq -1.7$	-0.897	0.315	32
TOA	$4.61 \leq \eta \leq 4.92$	3.75	0.0038	12
TOC	$-3.28 \leq \eta \leq -2.97$	-0.727	0.0038	12

Tab. 3.2.: Main details about the ALICE ITS [131].

Parameter	SPD	SDD	SSD
Total number of modules	240	260	1770
Number of readout channels per module	240	2×256	2×768
Radius (cm)	3.9 - 7.6	15.0 - 32.9	38.0 - 43.0
Material Budget per layer ($\%X_0$)	1.14 - 1.14	1.13 - 1.26	0.83 - 0.86
Spatial resolution $r\phi$ (μm)	12	35	20
Spatial resolution z (μm)	100	25	830
Two-track resolution $r\phi$ (μm)	100	200	300
Two-track resolution z (μm)	100	600	2400
Active cell size (μm^2)	50×425	202×294	95×40000

The low material budget and the high space resolution of the ITS allow for the reconstruction of primary¹ and secondary² vertices with a resolution better than 100 μm . In addition the ITS is used to extend the tracking down to $p_T = 80 \text{ MeV}/c$. The information about the resolution and the material budget of the ITS can be found in Tab. 3.2.

While the SPD is characterised by a digital readout, the readout of the SDD and of the SSD is analogical. For this reason the SDD and the SSD can be used for PID (Particle Identification) through the measurement of the specific energy loss dE/dx in the non relativistic region, below 200 MeV/c .

3.2.2 Time Projection Chamber

The TPC is the main tracking detector of the central barrel. It is optimised to provide momentum measurements with good two-track separation, particle identification via dE/dx measurement and vertex determination together with the other detectors of the central barrel. The TPC covers the pseudorapidity range $|\eta| < 0.9$ for tracks with full radial length, i.e. tracks with reconstructed points in ITS, TRD and TOF detectors. It provides measurements from low p_T (0.1 GeV/c) to high p_T (100 GeV/c) with a resolution on the transverse momentum of about 5% for $p_T < 20 \text{ GeV}/c$ and 35% at $p_T = 100 \text{ GeV}/c$ in standalone mode. The resolution improves of a factor of 5-7 (depending on the p_T) when combined with the ITS and the TRD information [110].

The TPC layout is shown in Fig. 3.3b: the detector consists of a cylindrical field cage, with a central high-voltage cathode, which divides the TPC in two halves, and

¹Primary vertex is a synonym of collision point.

²A secondary vertex (or decay vertex) is the point where an unstable particle decays.

two end-plate anodes equipped with readout detectors, consisting of multi-wire proportional chambers (MWPC). In each half of the TPC the central cathode and the end-plate anodes generate a highly uniform electric fields of 400 V/m. The total active volume has an inner radius of 85 cm, an outer radius of 250 cm and a total length along the beam direction of 500 cm, filled with a gas mixture. Two mixtures of gas were used: a mixture of Neon and CO₂ in Run 1 and in part of Run 2 and a mixture of Argon and CO₂ in the other part of Run 2.

When a charged particle passes through the active volume of the TPC it interacts with the gas, ionising it. The electrons produced in the ionisation drift towards the readout chambers in the end plates, providing the position in the transverse plane, while the position along z is provided by the measure of the drift time, for which it is necessary to use a trigger information from an external detector, like the T0. The position resolution is different along the radial direction because of the different track density: for the inner/outer radius is 1100/800 μm in the transverse plane and 1250/1100 μm along the beam axis. The TPC is also a PID detector: particle identification is performed via the measurement of the specific energy loss dE/dx , proportional to the number of collected electrons in the readout chambers. More details about the particle identification are given in section 3.4.

The TPC is an excellent detector for tracking and PID, but its cost is a large drift time (88 μs), which is a great limitation for the data acquisition rate (1 kHz).

3.2.3 Time Of Flight

The TOF detector consists of a large-area array of Multi-gap Resistive Plate Chambers (MRPC) in the central pseudorapidity range ($|\eta| < 0.9$). It is used for the PID in the intermediate momentum range, up to 2.5 GeV/ c for pions and kaons, up to 5 GeV/ c for protons and deuterons. The geometrical information about the TOF can be found in Tab. 3.1.

The PID is obtained through the measurement of the time of flight of a particle, from the collision point to the TOF detector. The collision time can be provided either by the T0 detector or via a combinatorial analysis of the particle arrival times at the TOF detector, as described in [132]. The overall TOF resolution, including the uncertainty on the start time of the event, is around 80 ps in Pb–Pb collisions and 120 ps in pp collisions. If the T0 information is not available, the bunch-crossing time is used and the resolution is about 200 ps. The complete review of the methods used to compute the event start time can be found in [132]. Combining the information

on the track length, the momentum and the time of flight it is possible to determine the mass of the particle through the relations, as described in Sec. 3.4.2.

3.2.4 VZERO

The V0 consists of two arrays of scintillator counters (VOA³ and VOC⁴). The geometrical information about the V0 can be found in Tab. 3.1. Both the detectors cover the full azimuthal angle interval. The time information provided by the V0 detector is used for the offline rejection of the machine-induced background (MIB), a particle flux coming from the interaction of the beam with the matter in the machine, typically the residual gas in the beam pipe. The rejection is based on the time difference between the signal in VOA and VOC. Indeed, this time difference for a collision is distributed around a known value (8.3 - 14 ns) and for MIB it is outside this region [113]. The V0 provides the Minimum-Bias trigger, defined as a logical or of VOA, VOC and SPD, and is used to measure the event multiplicity at forward rapidity (see Sec. 3.5), defined in terms of the joined signal of VOA and VOC (VOM).

3.2.5 TZERO

The detector consists of two arrays (T0A and T0C) of Cherenkov counters. The geometrical information about the T0 can be found in Tab. 3.1. Its main purpose is the determination of the event time (T0), independently from the track and vertex reconstruction, with a resolution below 50 ps. It can also be used for an independent determination of the vertex position along the beam line, with a precision of 1.5 cm, and to provide the first level (L0) trigger when the position is within the expected value.

3.3 Event reconstruction

The event reconstruction, both for real data and MC simulations, can be divided into several steps.

The first step is the local reconstruction: for each sub-detector the signal must be converted into the space-time information of the particle that generated it, forming a *recpoint* or *cluster*. Other information than the space-time coordinates, such as the PID information obtained via dE/dx , time of flight, etc., can be attached to the

³A = ATLAS: it is on the side towards the ATLAS experiment

⁴C = CMS: it is on the side towards the CMS experiment

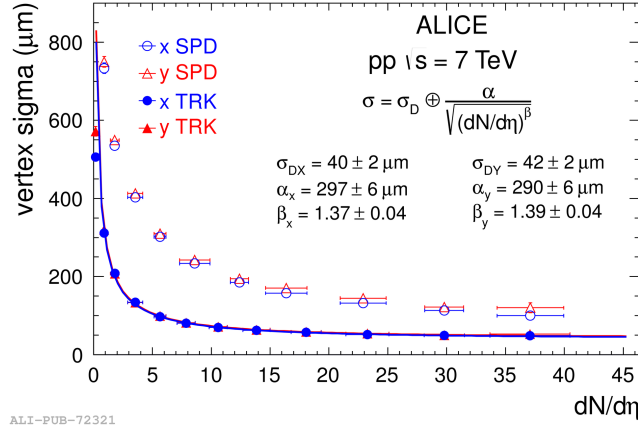


Fig. 3.4.: Resolution on the primary vertex position in pp collision at $\sqrt{s} = 7$ TeV, obtained with SPD-tracklets (SPD) and full tracks (TRK) [133].

recpoint.. Once the recpoints are available, it is possible to reconstruct the primary vertex and all the tracks.

The primary vertex can be reconstructed by using the SPD or the full reconstructed tracks. The second method is characterised by a higher resolution but, by definition, it can be performed only after the track reconstruction. The SPD reconstruction works as follows: the algorithm connects two points (one on each layer of the SPD), which are in the same azimuthal acceptance window, to form a segment called *tracklet*. The primary vertex position is obtained by minimising the distance among all the tracklets and removing the outliers. For this method at least two tracklets are necessary, but this condition is not always satisfied in pp collisions, where often just one tracklet is found. In this case it is possible to compute the z position of the vertex by using the beam line position in the transverse plane as a constraint. Similarly, the vertex reconstructed with the second method is obtained by minimising the distance among the full tracks. The resolution on the primary vertex position using SPD tracklets and full tracks is shown in Fig. 3.4. The resolution depends on the charged particle multiplicity: in Pb-Pb collisions, characterised by a higher multiplicity than pp collisions, a resolution of 10 μm is obtained, to be compared with the average resolution of 150 μm in pp collisions [133].

A schematic view of the track reconstruction is shown in Fig. 3.5. Tracks are reconstructed in the ITS and in the TPC with the Kalman filter algorithm [135], prolonged to the other detectors of the central barrel and then fitted again. The first track seeds are built starting from clusters in the two outermost pads of the TPC and the primary vertex reconstructed with the SPD tracklets. The seeds are then propagated inwards and the tracks are updated with every closest cluster found by the Kalman filter algorithm, until the inner radius of the TPC is reached. The efficiency of the reconstructed tracks in the TPC is shown in Fig. 3.6a as a function of p_T in pp collisions at $\sqrt{s} = 8$ TeV and in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV.

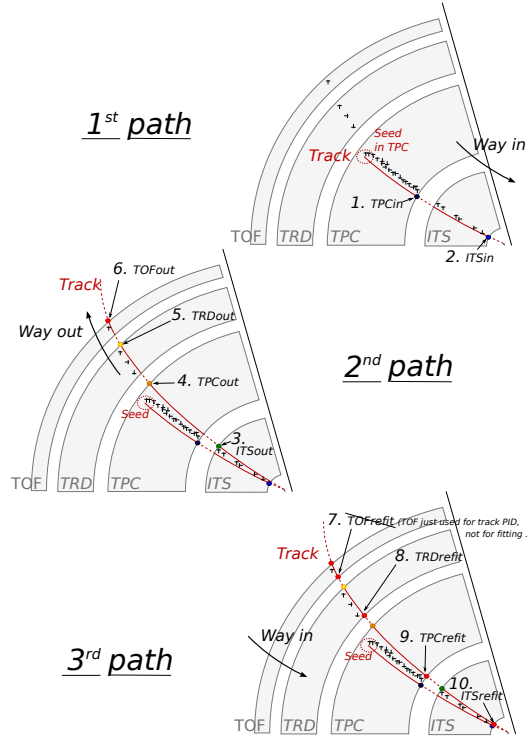


Fig. 3.5.: Schematic view of the three tracking steps of performed in the reconstruction of an event in ALICE. Fig. taken from Ref. [134].

The efficiency drops for $p_T \lesssim 500 \text{ MeV}/c$ due to the energy loss in the material of the detector, since the trajectory of particles that loose energy deviates from the helicoidal path of a charged track in the magnetic field. The observed shape of the efficiency at higher p_T is due to the loss of cluster in the dead zones between the TPC sectors.

The tracks reconstructed in the TPC are matched to the hits in the outermost SSD layer and become the seed for the track reconstruction in the ITS. Similarly to the track reconstruction in the TPC, the seed is propagated inwards and updated at each ITS layer, attaching all the clusters within a proximity cut. For each TPC track, a corresponding tree of track hypotheses in the ITS is produced. The candidates are selected according to their χ^2 . Fig. 3.6b shows the comparison between the fraction of tracks in TPC that have a prolongation in the ITS (matching efficiency), in real and simulated Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76 \text{ TeV}$, for different requirements on the ITS points [113].

In order to increase the efficiency for tracks with low transverse momentum ($p_T \leq 200 \text{ MeV}/c$) and to reconstruct the tracks of the particles traversing dead zones of the TPC or which decay before entering the TPC, the ITS clusters not attached to tracks propagated from the TPC are used to perform an ITS stand-alone reconstruction.

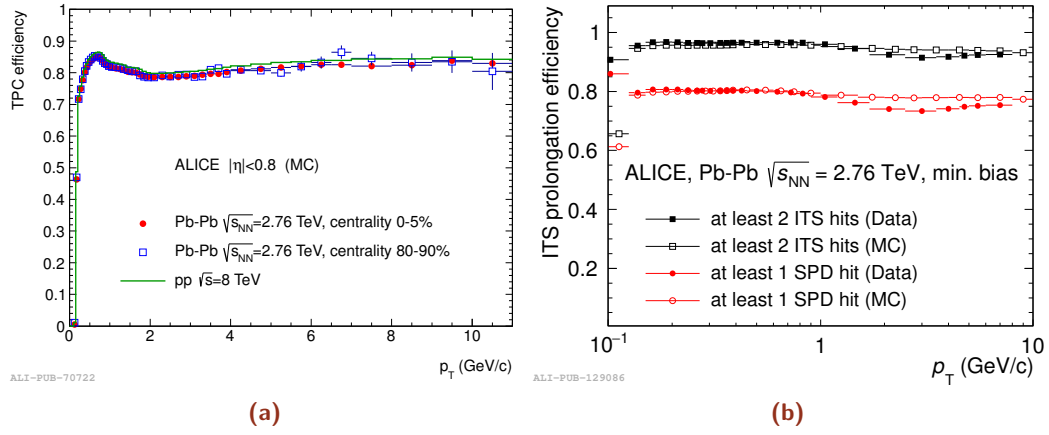


Fig. 3.6.: Reconstruction efficiency for TPC tracks in pp collisions at $\sqrt{s} = 8$ TeV and in central and peripheral Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV (a) and matching efficiency of tracks from ITS or TPC in real (full markers) and simulated (open markers) Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. Figures taken from Ref. [113].

The tracks obtained with clusters in both the ITS and the TPC detectors are then back-propagated outwards to match the clusters in the TRD and in the TOF and the signals in EMCAL, PHOS and HMPID. Finally, the tracks are re-fitted inwards using all the previously found clusters and propagated to the point of closest approach to the beam line. After the full tracking procedure, if two or more tracks are reconstructed, the position of the primary vertex is computed using global tracks. As shown in Fig. 3.4, the resolution on the primary vertex position is always higher when computed with full tracks than with SPD tracklets.

Fig. 3.7 shows the resolution of $1/p_T$ for TPC stand-alone tracks and tracks with both TPC and ITS clusters, with and without vertex constrain, in p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The resolution on $1/p_T$ is directly connected to that on the p_T by the relation:

$$\frac{\sigma_{p_T}}{p_T} = \frac{\sigma_{1/p_T}}{1/p_T}. \quad (3.1)$$

The resolution is quoted as $1/p_T$ because this can be extracted directly from the covariance matrix of the Kalman filter fit. The resolution on the reconstructed tracks vary from 1% to 10% in the p_T range between 0.1 and 100 GeV/c.

3.4 Particle identification

Five detectors are used for PID in different ranges of momentum: the ITS, the TPC, the TRD, the TOF and the HMPID. As already anticipated, the ITS and the TPC

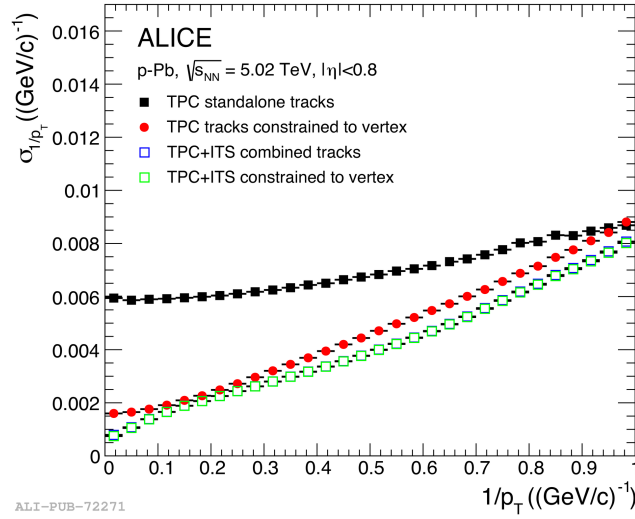


Fig. 3.7.: Resolution on the $1/p_T$ as a function of $1/p_T$ in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The resolution is reported for TPC standalone tracks, both with (red points) and without (black points) primary vertex constraint and for ITS+TPC tracks, with (green) and without (blue) vertex constraint. Figure taken from Ref. [113].

identify particles through the measurement of their specific energy loss, TRD via transition-radiation, TOF with the time of flight and finally HMPID via $\beta = v/c$, obtained from the Cherenkov angle measurement. The TPC and TOF are the main PID detectors used in this thesis and for this reason a dedicated description will be reported in the following for each of them.

3.4.1 TPC particle identification

Fig. 3.8 shows the TPC energy loss as a function of momentum. The expected detector response can be parameterised with splines or using a Bethe-Bloch parameterisation:

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

where β and γ are the relativistic factors and P_{1-5} are free parameters. This parameterisation can be fitted to the data and it is derived by the ALEPH collaboration [136].

The detector response to light hadrons can alternatively be parameterised using splines, provided by the central ALICE framework, as in Fig. 3.8. At low momenta ($p \leq 1 \text{ GeV}/c$) the dE/dx has a $1/\beta^2$ dependence on the momentum and the particle identification is performed with a fiducial selection around the expected signal. The width of the fiducial interval is usually expressed in terms of number of σ , being σ the expected dE/dx resolution for the analysed track. The resolution on the dE/dx is about 5.2% in pp collisions and 6% in central Pb–Pb collisions. At higher momenta

the signal extraction is still possible to extract the relative contributions of different particle species through the statistical unfolding of the dE/dx distributions. More details can be found in [110].

Since in Pb–Pb collisions the production yields of hadrons are higher than in pp collisions, with TPC and TOF it is possible to observe light (anti-)nuclei up to (anti-) ^4He . The panel in Fig. 3.8 shows the squared-mass distribution for $^3\bar{\text{He}}$ and $^4\bar{\text{He}}$ measured with the TOF. Using the TPC, ^3He and ^4He are easily separated from the other hadron species because the specific energy loss is proportional to the squared charge of the particle and their charge is the double of the charge of the other species. Using only the TPC it is possible to separate deuterons up to $p \sim 1.4 \text{ GeV}/c$. For larger momentum, the TOF is necessary.

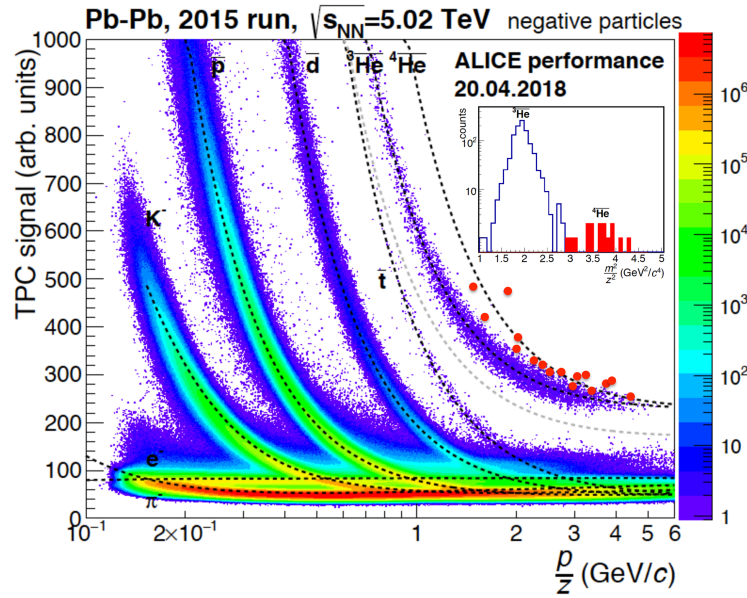


Fig. 3.8.: Specific energy loss of particles traversing the TPC volume as a function of their momentum in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 5 \text{ TeV}$. The dashed black lines represent the expected response of the TPC for different particle species. The panel in the top-right corner shows the squared-mass distribution for $^3\bar{\text{He}}$ and $^4\bar{\text{He}}$ measured with the TOF.

3.4.2 TOF particle identification

As anticipated in Sec. 3.2.3 the TOF is a PID detector that identifies particles by combining the information on the track length L , the momentum p and the time of flight t_{TOF} . Fig. 3.9 shows the relativistic β , measured with TOF, as a function of momentum in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02 \text{ TeV}$. The relativistic β is computed as:

$$\beta = \frac{L}{ct_{\text{TOF}}}. \quad (3.2)$$

In the figure one can distinguish the population of different particle species over a mismatch background, coming from the wrong association of a TOF point to a track. The mismatch background is more significant at high multiplicity, such as in Pb-Pb collisions.

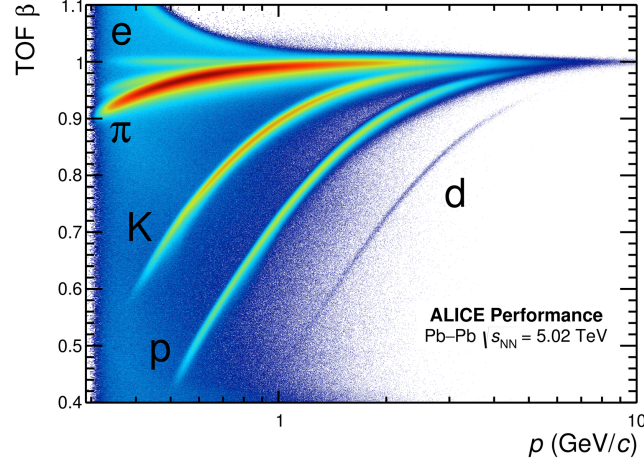


Fig. 3.9.: Relativistic β of the particles as a function of their momentum in Pb-Pb collisions at $\sqrt{s_{NN}} = 5$ TeV measured with the TOF.

3.5 Multiplicity determination

The charged particle multiplicity $dN_{ch}/d\eta$ is a key quantity for the event classification. Indeed, a multiplicity-based event classification permits the direct comparison between the results obtained in different collision systems. The charged particle multiplicity is determined from the VOM signal, the joint signal of V0A and V0C, which is proportional to the number of charged particles that pass through the detector. The

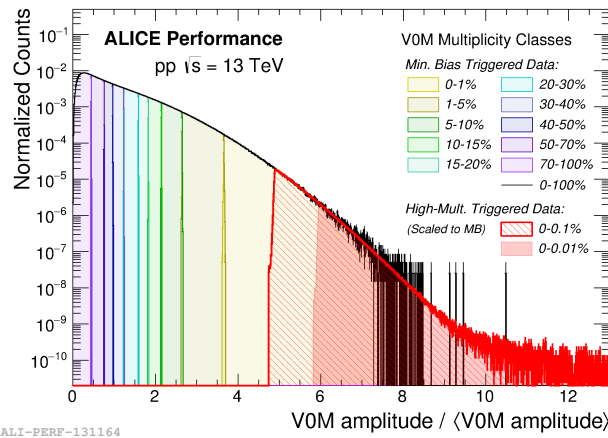


Fig. 3.10.: Distribution of the VOM amplitude/ $\langle$ VOM amplitude $\rangle$ ratio for pp collisions at $\sqrt{s} = 13$ TeV.

data sample is divided into ten multiplicity classes, defined from the percentiles of the distribution of the VOM signal amplitude, shown in Fig. 3.10. Each multiplicity class is identified by a roman number, going from I for the highest multiplicity class to X for the lowest multiplicity class. For example, the multiplicity class I includes all the top 1% multiplicity events, class II all the events ranking as 1-5% top multiplicity and so on down to the class X, that includes the events with a multiplicity in the range 70-100%.

The charged particle multiplicity at mid rapidity is estimated from the VOM signal through the comparison with a MC simulation: for each simulated collision, the simulated response of the V0 detector is related to the $dN_{\text{ch}}/d\eta$ at mid rapidity. In this way it is possible to get the correspondence between multiplicity and VOM signal.

3.6 Monte Carlo simulations

MC simulations are crucial to perform an analysis. They are used to evaluate the reconstruction efficiencies, to choose the criteria of track selection that maximise the signal extraction and to evaluate the systematic uncertainties. The first step is the event generation, i.e. the simulation of the interaction between two particles of the LHC beams in the interaction region. Simulated events are generated with Monte Carlo generators, such as PYTHIA [137] for pp collisions and HIJING [138] for Pb–Pb collisions. The outcome of this first step is a set of kinematic parameters of both stable and weakly-decaying particles that will be transported. Strong decays are handled in the generation step, due to their extremely low life-times (the mean lifetime of the ϕ meson is $\tau \approx 1.5 \times 10^{-22}$ s [4]).

The kinematic parameters are propagated through a detailed geometrical description of the ALICE detector using a transport framework. Three different transport codes are used in ALICE: GEANT3 [139], GEANT4 [140] and FLUKA [141]. These codes deal with the energy loss of particles in the detector material, the generation of secondary particles from the material and the decay of unstable particles.

The impact point of a particle on the sensitive area of a detector and its energy deposit form a *hit*. Hits are stored and then processed by the simulation of each detector response, to produce a signal in the detector electronic. Finally, the simulated signal is stored in the same raw format used in real data taking and reconstruction.

The details about the MC simulations used for this analysis will be given in Sec. 4.2.

Analysis techniques for the measurement of the (anti-)deuteron production spectra

The measurement of particle production spectra is a long procedure that consists of many analysis steps, going from the data selection to the final results. In this chapter the analysis techniques for the measurements of the (anti-)deuteron production spectra in pp collisions at $\sqrt{s} = 13$ TeV will be described. A general explanation of the analysis strategy will precede a detailed description of each analysis step.

4.1 Analysis strategy

The aim of this work is the measurement of (anti-)deuteron production spectra in pp collisions at $\sqrt{s} = 13$ TeV. The analysis has been carried out on the pp data samples collected in 2016 and 2017.

The first step of the analysis is the selection of good-quality pp collisions, from here on called *events*. For example, in order to maximise the number of particles in the detector acceptance, only those events with a reconstructed primary vertex close to the nominal interaction point, i.e. the origin of the reference frame as described in Sec. 3.2, are selected. This selection is performed by rejecting all the events whose distance between the primary vertex and the nominal collision point is larger than a value defined a priori, as reported in Sec. 4.3. Another important factor in the event selection is the rejection of multiple collisions that take place within the time resolution of the detector (300 ns for the SPD [131]). Such condition is called *pile-up*. There are two kinds of pile-up. *In-banch* pile-up takes place when multiple collisions occur in the same bunch crossing¹. *Out-of-banch* pile-up takes place when one or more collisions occur in bunch crossings that are different from the one that triggered the acquisition. While in-bunch pile-up is seen by all the detectors, out-of-bunch pile-up affects the different detectors according to their readout time. For example, the SPD has a readout time of 300 ns, while the TPC has

¹Bunch crossing is defined as the instant when the two LHC beams cross at the ALICE interaction point. Bunches are separated by time windows that are multiples of 25 ns, which is the period of the LHC clock.

a readout time of 100 μs [131]. Pile-up brings a bias in the determination of the number of events, affecting the normalisation of the production spectra, and a bias in the estimation of the charged particle multiplicity. For this reason, it is crucial to reject pile-up events. As it will be described in Sec. 4.3, an event selection based on multiple reconstructed vertices can remove both in-bunch and out-of-bunch pile-up. In addition, out-of-bunch pile-up determines a spatial shift of the reconstructed points in the drift detectors (SDD and TPC). In this case, pile-up is also removed with track selections, by requiring a match between tracks in the ITS and in the TPC, at least one reconstructed point in the SPD and the pointing to the primary vertex (see Sec. 4.4).

The second step is the selection of the reconstructed tracks. The selection criteria concern track-quality parameters (more details will be given in Sec. 4.4) and particle identification. Tracks are selected at mid-rapidity in the rapidity interval $|y| < 0.5$. In this work particle identification is performed using the TPC and the TOF detectors. The former is used for transverse momentum values lower than 1 GeV/c, because in this p_T region (anti-)deuterons can be easily separated from other particle species. For higher p_T values, (anti-)deuteron are identified through the time information provided by the TOF detector.

The number of (anti-)deuterons is evaluated in p_T intervals by counting the number of reconstructed tracks, with the appropriate p_T and positively identified as (anti-)deuterons. This operation is called *signal extraction*.

The results of the previous steps are the raw production spectra $d^2N_{raw}/dp_T dy$. To obtain the real production spectra some corrections must be applied:

$$\frac{d^2N}{dp_T dy} = \frac{d^2N_{raw}}{dp_T dy} \times \frac{1}{\epsilon \times Acc} \times f_{primary}. \quad (4.1)$$

The first correction factor of Eq. 4.1, $\epsilon \times Acc$, takes into account the detector efficiency and acceptance. Indeed, (anti-)deuterons in acceptance can possibly be lost because of inefficiencies related to the detectors, to the track reconstruction and to the particle identification. $\epsilon \times Acc$ is the ratio between the reconstructed tracks of (anti-)deuterons and all the (anti-)deuterons produced in a collision. Of course this information is not available in the experimental data and must be obtained from MC simulations. Indeed, in a simulated collision the number of generated (anti-)deuterons is known exactly, as well as the fraction that is correctly detected. The details about the computation of the $\epsilon \times Acc$ corrections are reported in Sec 4.6.

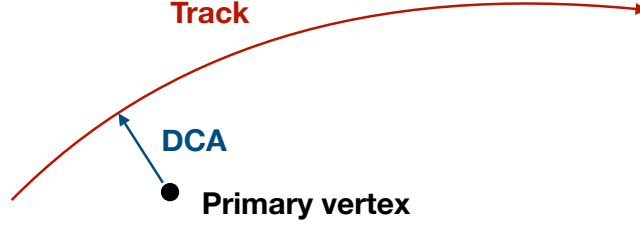


Fig. 4.1.: Definition of Distance of Closest Approach (DCA).

The second correction factor of Eq. 4.1 is the fraction of primary deuterons $f_{primary}$. To obtain the number of deuterons produced in the collision, we have to correct for the number of deuterons produced in secondary processes, the *secondary deuterons*. For example, part of the deuterons comes from the interaction of a neutron or a proton with the material of the detector, i.e. via material knock-out. Anti-deuterons cannot be produced via material knock-out as a consequence of baryon number conservation. Indeed, ordinary matter is composed of neutrons and protons, with baryon number $B = 1$, and the formation of an anti-deuteron, with baryon number $B = -2$, has an extremely high energy threshold and therefore it does not take place. Another potential source of secondary deuterons is the decay of hypertriton in its three-body decay channel ${}^3_{\Lambda}\text{H} \rightarrow \text{d} + \text{p} + \pi^-$ and its charge conjugated ${}^3_{\bar{\Lambda}}\bar{\text{H}} \rightarrow {}^3\bar{\text{d}} + \bar{\text{p}} + \pi^+$ for anti-hypertriton. However, the production of (anti-)hypertriton has not yet been observed in pp collisions and its production yield is therefore lower than that of ${}^3\text{He}$, which is less than a thousandth of the deuteron production rate [82]. The separation of primary and secondary deuterons is based on the measurement of the distance of closest approach (DCA), defined as the distance between the particle track and the primary vertex (see Fig. 4.1). The DCA is decomposed into its longitudinal component DCA_z and its transverse component DCA_{xy} . Both DCA_z and DCA_{xy} have a sign that depends on the parameterisation of the reconstructed track: if the coordinate of the track in the point of closest approach is larger than the coordinate of the vertex it is positive, otherwise it is negative. Since primary deuterons are produced in the collision, they are produced at the primary vertex and therefore their distance of closest approach is close to zero. Secondary deuterons are generally produced without a preferential direction and they show a broad DCA distribution. The expected DCA distributions for primary and secondary deuterons have to be studied with simulated data sets. In Sec. 4.7 it will be shown how the DCA distribution of secondary and primary deuterons are used as templates to fit the experimental data and get the primary fraction $f_{primary}$.

Finally, there is an additional correction coming from the normalisation of the production spectra. Due to trigger inefficiency, some events are erroneously rejected (*event loss*) and the (anti-)deuterons produced in such events are lost too (*signal loss*).

Therefore, it is necessary to correct the production spectra for the event loss (ϵ_{event}) and the signal (ϵ_{part}) losses, so that the final spectra can be computed as:

$$\frac{1}{N_{events}^{true}} \times \frac{d^2 N_{part}^{true}}{dp_T dy} = \frac{\epsilon_{events}}{N_{events}^{selected}} \times \frac{d^2 N_{part}^{selected}}{dp_T dy} \times \frac{1}{\epsilon_{part}(p_T)} \quad (4.2)$$

where N_{events}^{true} is the real number of events, $N_{events}^{selected}$ is the number of events selected by the trigger, N_{part}^{true} is the number of (anti-)deuterons produced in all the events and $N_{part}^{selected}$ is the number of (anti-)deuterons measured in the selected events. The signal-loss correction is defined as:

$$\epsilon_{part}(p_T) = \frac{d^2 N_{part}^{selected}}{dp_T dy} / \frac{d^2 N_{part}^{true}}{dp_T dy} \quad (4.3)$$

Both the correction factors ϵ_{event} and ϵ_{part} are computed from a MC simulation, because it is the only way to estimate the number of lost events and all the (anti-)deuterons lost in those events. In this work two kinds of normalisations are used. In the *Minimum Bias* analysis (see Sec. 3.2.4) all the hadronic inelastic pp collisions (*INEL* events) are considered. In the analysis of (anti-)deuteron production as a function of multiplicity, the presence of at least a charged particle in the acceptance ($|\eta| < 1$) is required (*INEL* > 0). A detailed description of the computation of these correction factors will be given in Sec. 4.8.

4.2 Data and Monte Carlo samples

This work is based on the datasets of minimum-bias pp collisions at $\sqrt{s} = 13$ TeV collected in 2016 and 2017. The total number of MB events collected is 1.1 billion events.

Most of the corrections to be applied to the raw spectra are obtained from MC simulations. The MC samples used for this work were generated with the PYTHIA8-Monash2013 event generator [137], using GEANT3 [139] to simulate the transport of the particles that are stable for strong-interaction through the material. The strong decay of particles is carried out by the event generator. The simulation is set up to reproduce the detector configuration that occurred during data acquisition. Since with Pythia event generator (anti-)deuterons are not produced, a dedicated particle generator is used to inject particles on top of the simulated event. The transverse momentum of the injected particles is taken randomly from a flat distribution from 0 to 6 GeV/c, the azimuthal angle from a flat distribution in the range $[0, 2\pi]$ and the rapidity from a flat distribution in the range $|y| < 1$. On top of each event, one particle of those listed in Tab. 4.1 is injected. Only one particle per event is injected not to bias the charged-particle multiplicity of the corresponding event. This

Tab. 4.1.: Details about the injection scheme adopted for the MC samples used for this analysis.

MC injection scheme	
Rapidity range	$ y < 1$
Azimuthal range	$0 \leq \varphi < 2\pi$
p_T range	$0 \leq p_T < 6 \text{ GeV}/c$
Injected particles	$d, \bar{d}, t, \bar{t}, {}^3\text{He}, {}^3\bar{\text{He}}, {}^4\text{He}, {}^4\bar{\text{He}}$
Injection scheme	1 per event

MC sample consists of 106 million events. Since secondary deuterons are scarcely produced, other 550 million MC events without injected (anti-)nuclei are used to evaluate the DCA_{xy} distribution of secondary deuterons.

4.3 Event selection

After the online event selection described in Sec. 3.2.4, a further event selection is performed offline. The first offline event selection concerns the primary vertex reconstruction: only those events with a reconstructed primary vertex are selected. As mentioned in Sec. 3.3, the primary vertex can be reconstructed both using full tracks and using SPD tracklets and the resolution on its position is higher with full tracks than with SPD tracklets. It is sufficient that the primary primary is reconstructed in one of the two ways. If a primary vertex reconstructed with SPD tracklets is available, the resolution on its position along the beam direction must be better than 0.25 cm. If the vertex is reconstructed with both the methods, the distance along the beam direction between the two points must be shorter than 0.5 cm. Then, in order to have a good acceptance, all the events with a reconstructed primary vertex outside the region of ± 10 cm along the beam direction from the nominal collision point are rejected ($|V_z| \leq 10$ cm).

The next selection concerns pile-up rejection. There are two levels of pile-up rejection. The first level is based on the number of primary vertices reconstructed with the SPD with more than n contributors, where n is a parameter of the selection. The number of contributors is the number of SPD tracklets² used for the vertex reconstruction. Based on the multiplicity dependence of the false positive pile-up in previous analyses, the n parameter has been set to 3 for events with less than 20 tracklets, to 5 for events with more than 50 tracklets and to 4 for the remaining

²see definition in Sec. 3.3

Tab. 4.2.: Summary of the selections of the events used in the analysis. The variables are described in the text.

Event selections	
Variable	Selection
Trigger	MB
$ V_z $	≤ 10 cm
z resolution on SPD vertex	≤ 0.25 cm
Multiple SPD vertices	rejected with more than n contributors
$ \Delta V_z $	≤ 0.5 cm

events. This selection can remove the pile-up of collisions taking place in the same bunch crossing and of the collisions out of bunch within the SPD readout time (300 ns). However, with this method it is not possible to resolve collisions closer than 8 mm along the beam direction. In this case pile-up is not detected and the two collisions are treated as a single one. The second level of pile-up rejection concerns the distribution of the z coordinate of the primary vertices, in particular the difference between the z coordinate of the primary vertex obtained from the SPD based method and that obtained with the track based method (ΔV_z). For this reason, the events in which the difference between the z coordinates obtained with the two different methods is larger than 0.5 cm are rejected ($|\Delta V_z| < 0.5$ cm).

Finally, in the analysis of the (anti-)deuteron production as a function of multiplicity, INEL > 0 events are selected by requiring offline at least one charged particle in the pseudo-rapidity region $|\eta| < 1$, on top of the minimum bias trigger. A summary of all the event selections is reported in Tab. 4.2. After all the selections, 970 million events survive out of 1.1 billion minimum bias events. Fig. 4.2 shows the V0M multiplicity percentile distribution of the selected events (see Sec. 3.5). The distribution is flat, as expected for a data sample collected with a minimum-bias trigger, where no particular selection on the event multiplicity is required. Based on the available data, the sample has been divided into nine multiplicity classes, as reported in Tab. 4.3. The number of selected events (970 million) is much larger than that used for the analysis about deuteron production in pp collisions at $\sqrt{s} = 7$ TeV (240 million) [85], where the data sample was divided into five multiplicity classes. Moreover, it has been possible to measure (anti-)deuteron production in pp collisions up to the same charged particle multiplicity ($dN_{\text{ch}}/d\eta \sim 26$) as measured in the high multiplicity p–Pb collisions at similar centre-of-mass energies.

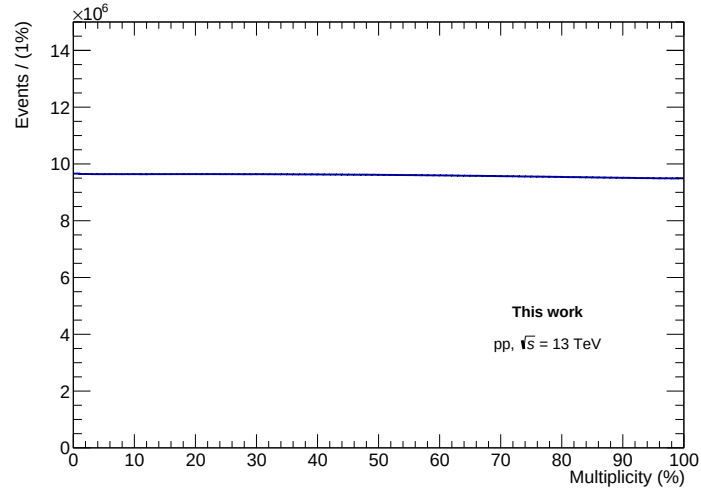


Fig. 4.2.: VOM multiplicity percentile distribution of the analysed data sample after the event selection. The VOM multiplicity percentile increases when the multiplicity of the selected event decreases.

Tab. 4.3.: Multiplicity classes used in this analysis, expressed in terms of percentiles of the V0 signal distribution (see Sec. 3.5). The last column contains the average charged particle multiplicity at mid rapidity corresponding to each class.

Multiplicity class	Multiplicity percentile (%)	$\langle dN_{\text{ch}}/d\eta \rangle$
I	0 - 1	26.02 ± 0.35
II	1 - 5	20.02 ± 0.27
III	5 - 10	16.17 ± 0.22
IV + V	10 - 20	12.91 ± 0.13
VI	20 - 30	10.02 ± 0.14
VII	30 - 40	7.95 ± 0.11
VIII	40 - 50	6.32 ± 0.09
IX	50 - 70	4.50 ± 0.07
X	70 - 100	2.55 ± 0.04

4.4 Track selection

In this analysis only the tracks in the mid-rapidity region are considered ($|y| < 0.5$). Moreover, only the tracks in the pseudorapidity region $|\eta| < 0.8$ are selected, in order to use only the geometrical region in which the experiment provides full track reconstruction and the best PID information.

The selected tracks are required to have at least 70 clusters in TPC to guarantee a transverse momentum resolution of 2% in the p_T range considered and a dE/dx resolution of about 6%. Requiring two points in the ITS, out of which at least one in the SPD, it is possible to assure a resolution better than 300 μm [133] on the distance of closest approach to the primary vertex both in the plane perpendicular to the beam direction (DCA_{xy}) and along the beam direction (DCA_z).

With a selection of the tracks with $|\text{DCA}_z| \leq 1$ cm and $|\text{DCA}_{xy}| \leq 0.12$ cm it is possible to reduce the contribution of the secondary deuterons, since they do not generally point to the primary vertex and they show a broad DCA distribution. Primary deuterons are instead characterised by a DCA distribution peaked at zero.

In addition, only the tracks with χ^2 per ITS cluster smaller than 36 and χ^2 per TPC cluster smaller than 4 are selected. The selected deuterons are required to be refitted both in TPC and ITS and not to come from a *kink topology*, i.e. coming from a two-body weak decay together with a neutral particle. Finally, both a refit of the track in the ITS and in the TPC is requested (see Sec. 3.3). The track selections are summarised in Table 4.4.

4.5 Signal extraction

The selected tracks are separated into p_T bins according to their transverse momentum and for each p_T bin the signal extraction is performed by using the PID power provided by TPC and TOF. At low transverse momentum ($p_T < 1$ GeV/ c), the only TPC is sufficient to separate (anti-)deuterons from other species. For higher transverse momentum ($p_T > 1$ GeV/ c) TOF is used, after a 3σ selection with the TPC in order to reduce the background coming from misidentified particles and from mismatch, i.e. the incorrect association of a TOF hit with a reconstructed ITS+TPC track

Tab. 4.4.: Summary of the track selections used in this analysis.

Track selections	
Variable	Selection
$ \eta $	< 0.8
$ y $	< 0.5
$n_{TPCclusters}$	> 69
$n_{ITSclusters}$	> 1
$n_{SPDclusters}$	> 0
$ DCA_z $	$< 1 \text{ cm}$
$ DCA_{xy} $	$< 0.12 \text{ cm}$
$\chi^2_{TPC}/n_{TPCclusters}$	≤ 4
$\chi^2_{ITS}/n_{ITSclusters}$	≤ 36
kink topologies	rejected
ITS refit	yes
TPC refit	yes

4.5.1 (Anti-)deuteron identification with the TPC

The signal extraction with the TPC is based on the measurement of the specific energy loss in the TPC active volume dE/dx , in particular on the difference between the measured dE/dx and the expected one for the particle of interest, expressed in terms of number of sigmas $n\sigma$:

$$n\sigma = \frac{(dE/dx)_{measured} - (dE/dx)_{expected}^i}{\sigma^i} \quad (4.4)$$

where $(dE/dx)_{measured}$ is the measured specific energy loss of the track candidate, $(dE/dx)_{expected}^i$ is the expected specific energy loss for the particle species i (in this case a(n) (anti-)deuteron), and σ^i is the detector resolution for the species i . Measured dE/dx as a function of particle rigidity p/z is shown in Fig. 4.3.

Fig. 4.4 shows the fit to the $n\sigma$ distribution for deuterons in the multiplicity class I, corresponding to the highest multiplicity, for the p_T ranges $0.8 \leq p_T < 0.9 \text{ GeV}/c$

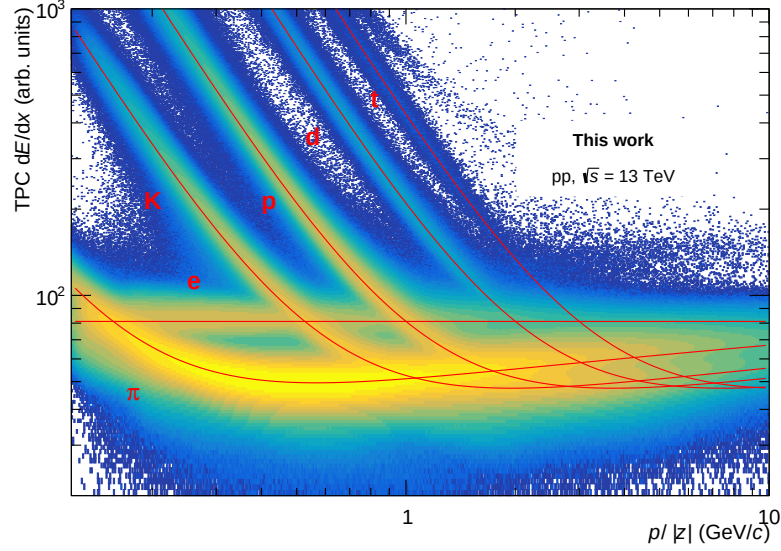


Fig. 4.3.: Specific energy loss in the TPC active volume as a function of particle rigidity in pp collisions at $\sqrt{s} = 13$ TeV. The solid red lines represent the expected response of the TPC for different particle species.

(Fig. 4.4a) and $0.9 \leq p_T < 1.0$ GeV/c (Fig. 4.4b). The function used for the signal is a gaussian with two exponential tails, defined as:

$$s(x; N_{sig}, \mu, \sigma, \alpha_0, \alpha_1) \propto N_{sig} \begin{cases} \exp[-\alpha_1 \left(\frac{x-\mu}{\sigma} - \frac{\alpha_1}{2} \right)] & \text{for } x < \mu + \alpha_1 \sigma \\ \exp[-\frac{1}{2} \left(\frac{x-\mu}{\sigma} \right)^2] & \text{for } \mu + \alpha_1 \sigma \leq x < \mu + \alpha_0 \sigma \\ \exp[-\alpha_0 \left(\frac{x-\mu}{\sigma} - \frac{\alpha_0}{2} \right)] & \text{for } x > \mu + \alpha_0 \sigma \end{cases} \quad (4.5)$$

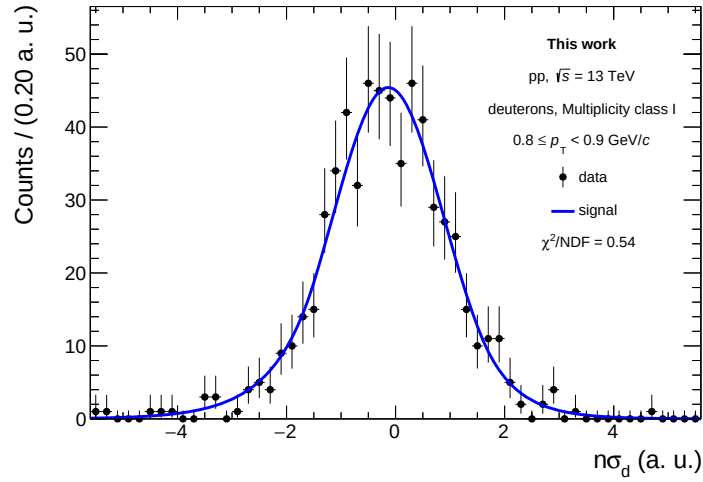
where N_{sig} is the number of signal counts, μ and σ are respectively the mean and the standard deviation of the Gaussian and α_0 (α_1) is the number of σ s from μ on the right (left) at which the function becomes an exponential. The slope of the exponential is defined so that the s function and all its derivatives are continuous and differentiable. The fit is performed maximising the extended likelihood, so that the number of (anti-)deuterons is directly obtained as a fit parameter (N_{sig}).

For $p_T > 0.9$ GeV/c, some background coming from other species is present. The background is fitted with an exponential with parameter τ_0 and with normalisation N_{bkg} , as shown in Fig. 4.4b. The signal function is always parameterised as a double-tailed Gaussian. By maximising the extended likelihood it is possible to get directly the number of count corresponding to background (N_{bkg}) and signal (N_{sig}).

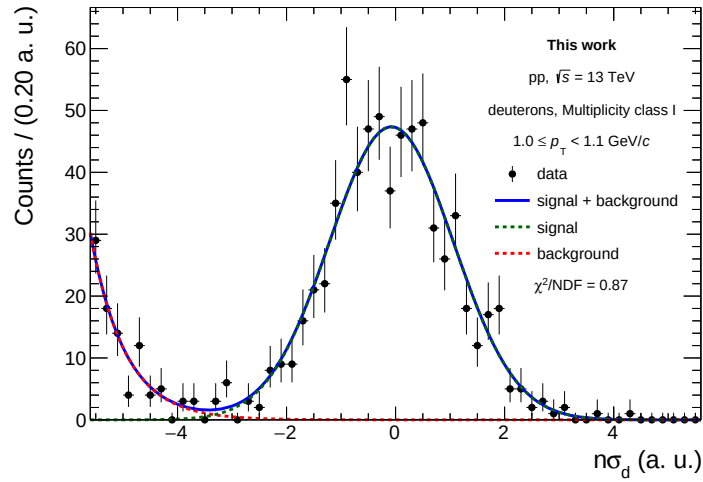
Signal extraction with with the only TPC is possible for $p_T < 1.2$ GeV/c, as shown in Fig. 4.5b. For $p_T \geq 1.2$ GeV/c the background becomes the dominant part. In Fig. 4.5, it is possible to appreciate how the background contribution becomes more

Tab. 4.5.: p_T intervals for the signal extraction with TPC.

p_T intervals (GeV/c)	Background
$[0.7, 0.8)$, $[0.8, 0.9)$	no
$[0.9, 1.0)$, $[1.0, 1.1)$, $[1.1, 1.2)$	exponential

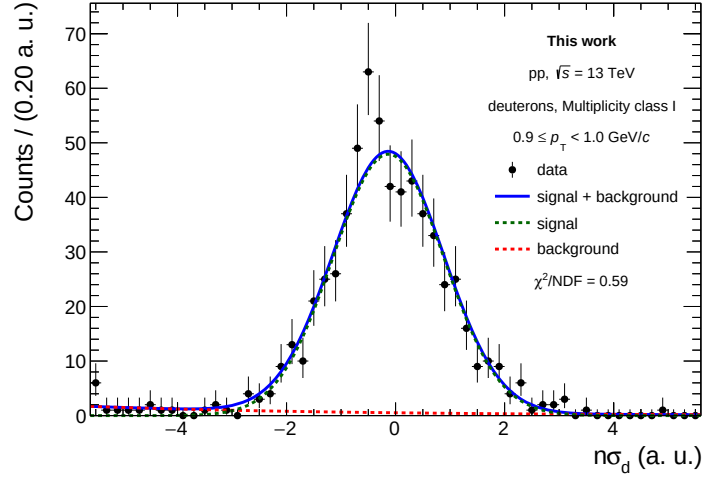


(a)

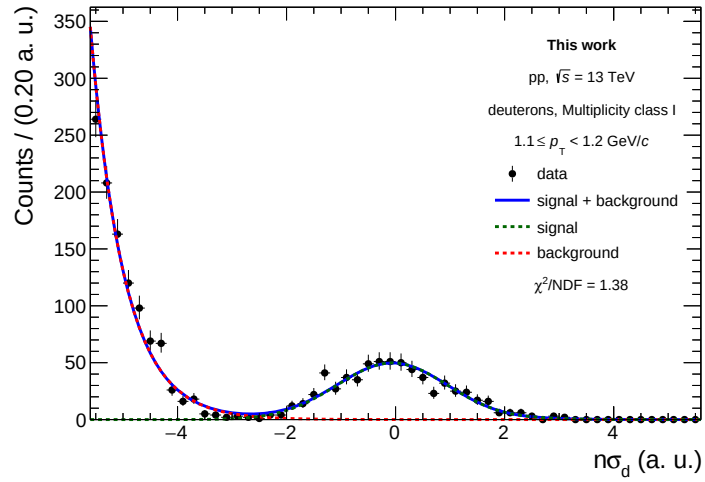


(b)

Fig. 4.4.: Fits to the TPC $n\sigma$ distributions for deuterons, in the multiplicity class I, in the p_T ranges $0.8 \leq p_T < 0.9$ GeV/c (a) and $1.0 \leq p_T < 1.1$ GeV/c (b). In (a) the background is negligible.



(a) Minimum background



(b) Maximum background

Fig. 4.5.: Fits to the TPC $n\sigma$ distributions of deuterons, in the multiplicity class I, for the p_T -intervals with minimum (top) and maximum background (bottom).

and more important with increasing transverse momentum. Raw yield extraction with TPC is performed in the p_T intervals reported in Tab. 4.5. However, even if signal extraction with the TPC is still possible, for $p_T \geq 1$ GeV/ c the signal is extracted using the TOF detector. The fits for all the p_T intervals in all the analysed multiplicity classes are omitted for a matter of space (100 plots) but they are very similar to those reported in Fig. 4.4 and Fig. 4.5.

4.5.2 (Anti-)deuteron identification with the TOF

The signal extraction with the TOF detector is based on the measurement of the squared mass m^2 of the particle associated to the track:

$$m^2 = \frac{p^2}{c^2} \left(\frac{1}{\beta^2} - 1 \right) \quad (4.6)$$

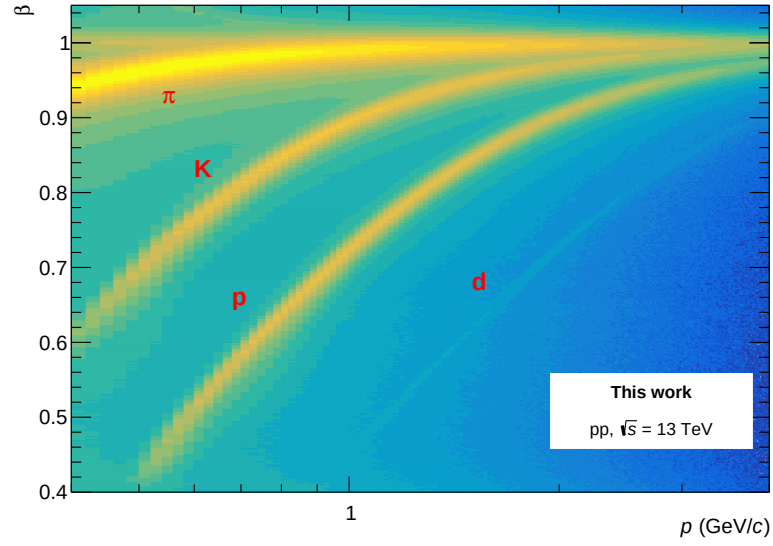
where p is the particle momentum and β is the particle velocity divided by the speed of light c . β is then computed as $\beta = L/(c t_{TOF})$, being L the length of the track and t_{TOF} the time of flight of the particle. In Fig. 4.6 the distribution of β as a function of momentum p is reported. In Fig. 4.6a it is possible to see how (anti-)deuterons, compared to the other particle species, are not well separated from the background coming from mismatch, i.e. the wrong association of a TOF hit with a track. For this reason, in the TOF analysis a pre-selection is performed, based on the TPC dE/dx : only the tracks with $|n\sigma_{TPC}| < 3$ are selected. As shown in Fig. 4.6b, in this way it is possible to reduce the background considerably. For simplicity, the m^2 distributions are centred on the deuteron squared-mass m_d^2 .

The (anti-)deuteron signal is modelled using a Gaussian with an exponential tail on the right:

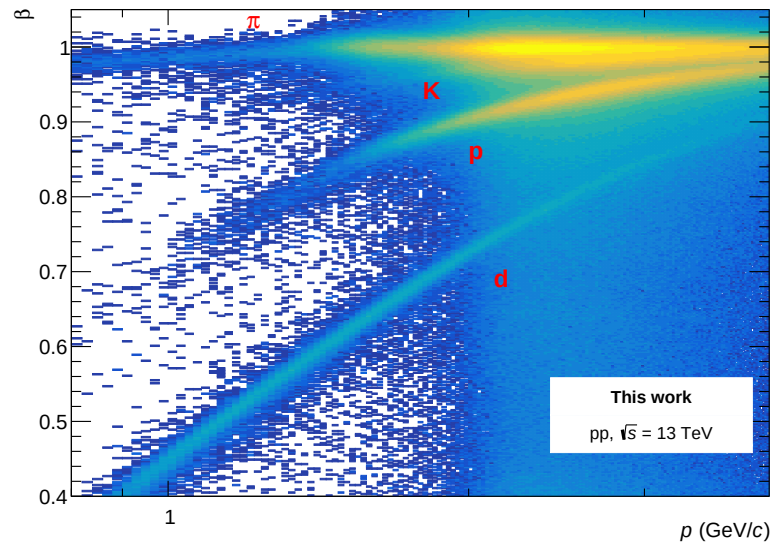
$$s(x; N_{sig}, \mu, \sigma, \alpha) \propto N_{sig} \begin{cases} \exp[-\frac{1}{2} \left(\frac{x-\mu}{\sigma} \right)^2] & \text{for } x \leq \mu + \alpha\sigma \\ \exp[-\alpha \left(\frac{x-\mu}{\sigma} - \frac{\alpha}{2} \right)] & \text{for } x > \mu + \alpha\sigma \end{cases} \quad (4.7)$$

where N_{sig} is the number of signal counts, μ and σ are respectively the mean and the standard deviation of the Gaussian and α is the number of standard deviations from μ at which the function becomes an exponential. As for the TPC, the slope of the exponential is defined so that the s function and all its derivatives are continuous and differentiable.

For $p_T < 1.8$ GeV/ c , the TPC pre-selection is sufficient to reject the background (see Fig. 4.7a). Therefore, in this p_T region the signal is extracted by counting the



(a) Without TPC pre-selection



(b) With TPC pre-selection

Fig. 4.6.: Measured TOF β as a function of momentum.

Tab. 4.6.: p_T intervals for the signal extraction with TOF.

p_T intervals (GeV/c)	Background
[0.9, 1.0) , [1.0, 1.1) , [1.2, 1.4) , [1.4, 1.6) , [1.6, 1.8)	no
[1.8, 2.0) , [2.0, 2.2) , [2.2, 2.6) , [2.6, 3.0) , [3.0, 3.4) , [3.0, 3.8)	double exponential

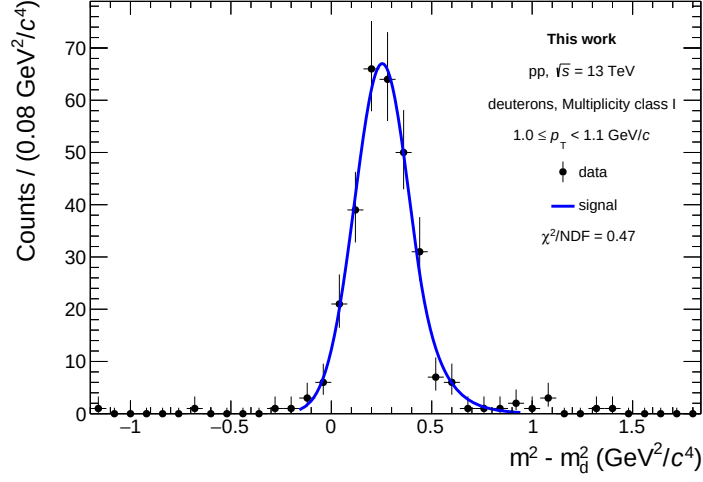
candidates. This procedure is called *bin counting*. First, the distribution is fitted with the function reported in Eq. 4.7. From this fit to the m^2 distribution μ and σ are obtained. Then, bin counting is performed in the region $\mu - 3\sigma < m^2 - m_d^2 < \mu + 5\sigma$. An asymmetric interval is chosen to take into account the right exponential tail. The position and the width of the interval are important because they are varied to determine the systematic uncertainty related to the signal extraction (see Sec. 4.9).

For $p_T > 1.8$ GeV/c background is present. Background is modelled as the sum of two exponentials:

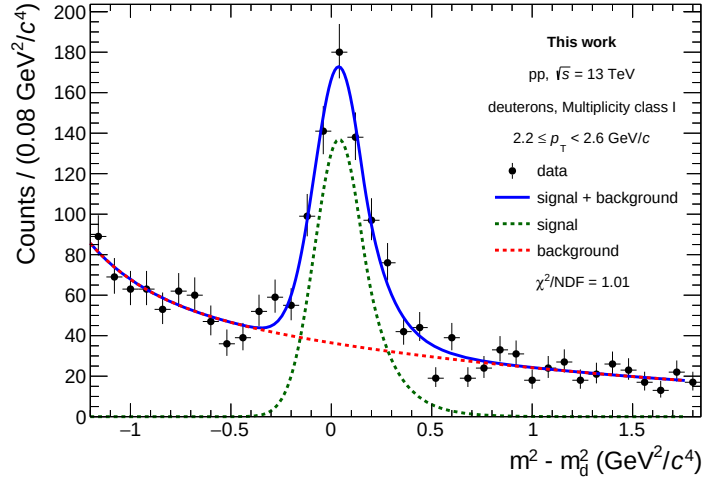
$$b(x; k, \tau_0, \tau_1) \propto k e^{-\tau_0 \cdot x} + (1 - k) e^{-\tau_1 \cdot x} \quad (4.8)$$

where τ_0 and τ_1 are the slopes of the two exponentials and k is the weight of those, limited between 0 and 1. In this p_T region where background is present, the m^2 distribution is fitted through the maximisation of the extended likelihood, so that the number of (anti-)deuterons N_{sig} is obtained as a fit parameter. In Fig. 4.7 it is possible to notice that the mean value of the m^2 distribution is not centred at 0 but shifted towards higher m^2 values and that this shift is more important at low p_T . This issue, already known from previous measurements in pp collisions at $\sqrt{s} = 7$ TeV [85], in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] and in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80], is due to the response of the TOF for long times of flight, i.e. small β . However, this behaviour is reproduced in MC simulations and it does not affect the signal extraction because μ is taken as a fit parameter. As for the TPC analyses, the fits for all the p_T intervals in all the analysed multiplicity classes are omitted for a matter of space (220 plots) but they are very similar to those reported in Fig. 4.7.

In Fig. 4.8 the raw spectra obtained from the signal extraction are showed for deuterons (Fig. 4.8a and Fig. 4.8b) and anti-deuterons (Fig. 4.8c and Fig. 4.8d). The spectra are divided by the number of events of the corresponding multiplicity class and scaled for a constant value for visibility. To obtain the final spectra from the raw ones, the content of each bin of the histogram must be divided by the bin width and the corrections described in Sec. 4.1 must be applied. The corrections are based on MC simulations and are described on the following sections.



(a)



(b)

Fig. 4.7.: Fits to the TOF squared-mass m^2 distributions for deuterons, in the multiplicity class I, in the p_T ranges $1.0 \leq p_T < 1.1$ GeV/c (top) and $2.2 \leq p_T < 2.6$ GeV/c (bottom). The two distributions are centred on the deuteron squared-mass m_d^2 . In (a) the background is negligible and the signal is extracted via bin counting.

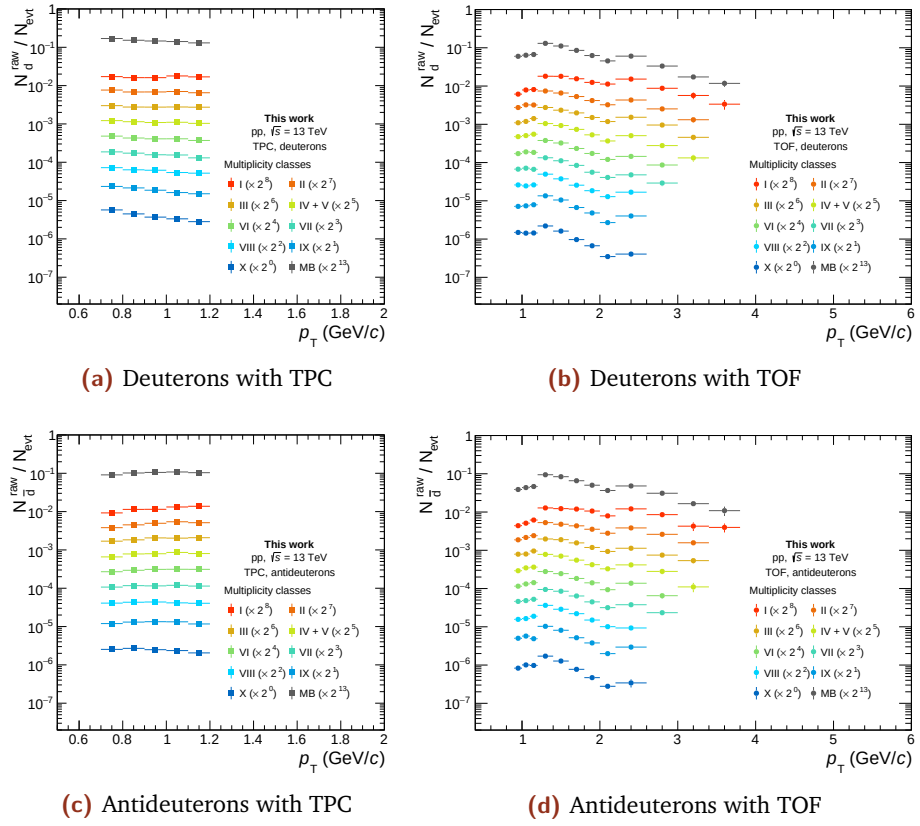


Fig. 4.8.: Raw spectra for deuterons ad anti-deuterons normalised by the number of events in the corresponding multiplicity class.

4.6 Correction for the detector efficiency and acceptance

As described in Sec. 4.1, not all the trajectories of the (anti-)deuterons produced in a collision are reconstructed. The fraction of generated (anti-)deuterons that are correctly reconstructed is evaluated from the MC sample. In particular, the correction that takes into account the detector efficiency and acceptance is defined as:

$$\epsilon \times Acc(p_T) = \frac{N_{rec}(|y| < 0.5, |\eta| < 0.8; p_T)}{N_{gen}(|y| < 0.5; p_T)}, \quad (4.9)$$

where N_{rec} and N_{gen} are the number of reconstructed and generated (anti-)deuterons, respectively. Both N_{rec} and N_{gen} are evaluated in the p_T intervals used for the signal extraction and reported in Tab. 4.5 and in Tab. 4.6. N_{gen} includes all the generated particles in the rapidity range $|y| < 0.5$, because the production spectra are measured in this rapidity range. N_{rec} , instead, includes all the reconstructed tracks in the rapidity range $|y| < 0.5$ and in the pseudo-rapidity range $|\eta| < 0.8$, because these are the values of the track selections, reported in Tab. 4.4.

In Fig. 4.9 $\epsilon \times Acc$ as a function of p_T is shown for the MB analysis. The correction is systematically higher for the TPC analysis (squares) with respect to the TOF analysis (circles). This almost 20% difference is due to the (anti-)deuteron absorption in the material between the TOF and the TPC. The absorption is mainly due to the presence of the TRD, which is characterised by an average radial thickness of 25% of a radiation length X_0 [117]. Moreover, $\epsilon \times Acc$ is consistently higher for deuterons than for anti-deuterons, reflecting a higher absorption of anti-deuterons. The low value of $\epsilon \times Acc$ for the TOF analysis in the transverse momentum region $p_T < 1$ GeV/c is the reason why the TOF analysis cannot be extended towards lower p_T , both because of the absorption and because at low p_T fewer particles are in the detector.

Finally, the correction has been computed separately for each multiplicity class and is compared with that calculated for the minimum bias sample. Fig. 4.10 shows that the ratio between $\epsilon \times Acc$ computed in multiplicity classes and with the MB sample is always compatible with unity within the uncertainties. For this reason the $\epsilon \times Acc$ correction computed for the MB analysis is used for all the multiplicity classes, since it is obtained from a larger data sample and hence shows smaller uncertainties.

4.6.1 Hadronic interaction correction

The MC samples used for the estimation of the detector efficiency are produced using GEANT3 [139] as transport code. The transport code is responsible for

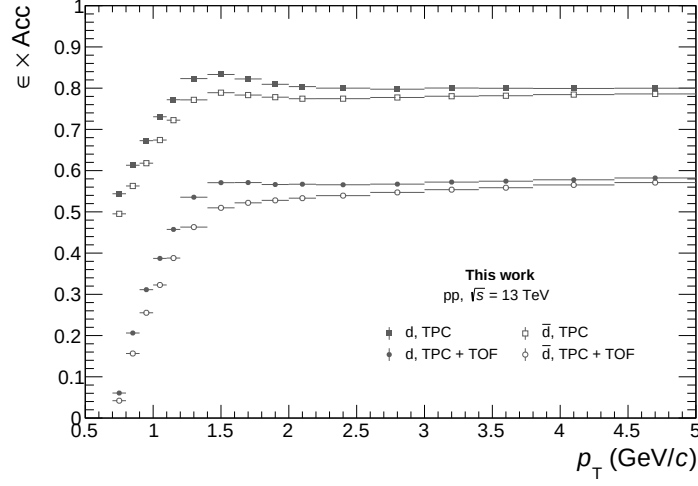
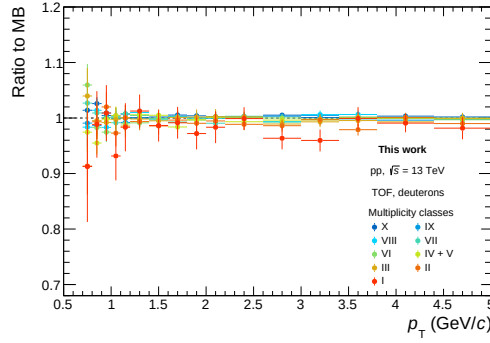
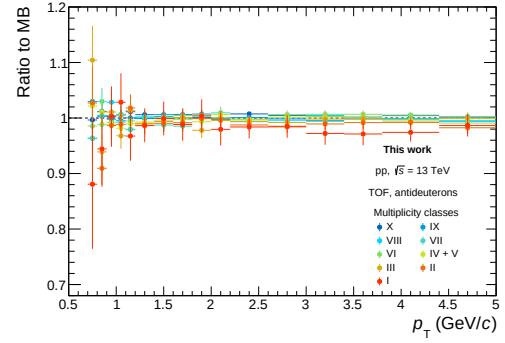


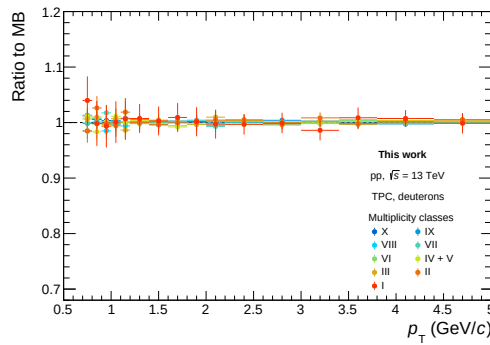
Fig. 4.9.: $\epsilon \times Acc$ correction as a function of transverse momentum for deuterons (full markers) and anti-deuterons (open markers), for TPC (squares) and TOF (circles), for the MB analysis.



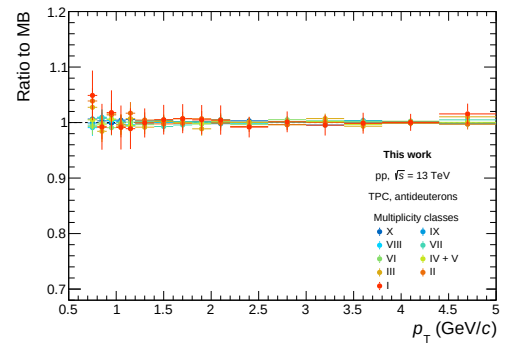
(a) Deuterons, TOF



(b) Antideuterons, TOF



(c) Deuterons, TPC



(d) Antideuterons, TPC

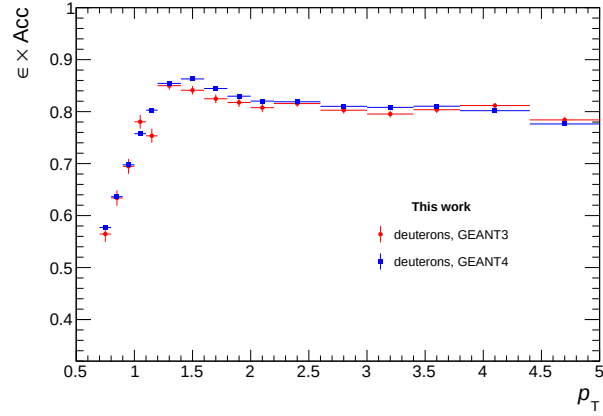
Fig. 4.10.: Ratio between the efficiencies computed in multiplicity classes and in the MB sample.

Tab. 4.7.: Details about the injection scheme adopted for the studies on the hadronic interaction correction.

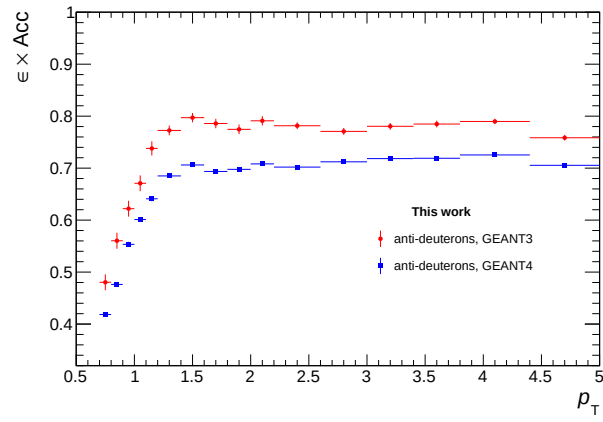
MC injection scheme	
Rapidity range	$ y < 0.8$
Azimuthal range	$0 \leq \varphi < 2\pi$
p_T range	$0 \leq p_T < 5 \text{ GeV/c}$
Injected particles	$p, \bar{p}, d, \bar{d}, t, \bar{t}, {}^3\text{He}, {}^3\bar{\text{He}}, {}^4\text{He}, {}^4\bar{\text{He}}$
Injection scheme	1 of each kind per event
Underlying event	none

the propagation of the particles produced by the event generator. It takes care of the interaction of the particles with the detector and of the weak decay of unstable particles. All these processes are determined by the definition, inside the transport code, of the cross sections that regulate the physics processes. However, for (anti-)deuterons this aspect is tricky because the hadronic interaction cross section at the energies considered in this analysis have not been measured. In ALICE simulations the standard transport code is GEANT3. However, it is known from other ALICE analyses on nuclei [80] that the absorption cross sections used in GEANT4 [140] provides a more realistic absorption of the nuclei and anti-nuclei.

In order to provide a better description of (anti-)nuclei absorption, the efficiencies computed with the standard transport code are scaled to reproduce those obtained with GEANT4. The correction has been estimated from two MC simulations in which only nuclei are generated, without any underlying collision. The adopted injection scheme is reported in Tab. 4.7. Each sample consists of about 50 million events and one is based on GEANT3 as transport code, the other on GEANT4. The efficiencies are computed for both the MC samples as described in Sec. 4.6. In Fig. 4.11, $\epsilon \times Acc$ has been evaluated for deuterons (Fig. 4.11a) and for anti-deuterons (Fig. 4.11b). In this case $\epsilon \times Acc$ is that evaluated for the TPC analysis. While the two transport codes does not show difference for deuterons, for anti-deuterons GEANT3 seems to underestimate the absorption with respect to GEANT4. In particular, the ratio between $\epsilon \times Acc$ evaluated with the two transport codes as a function of p_T is shown in Fig. 4.12. The dashed lines represent the mean values of the ratios for deuterons and anti-deuterons. While for deuterons the ratio is compatible with the unity, for anti-deuterons an average ratio of 1.1 is observed. For this reason, $\epsilon \times Acc$ is divided by 1.1 for anti-deuterons, both for the TPC and for the TOF analyses, because particles that reach the TOF detector have travelled all the path from the collision point to the outer wall of the TPC.



(a) deuterons



(b) anti-deuterons

Fig. 4.11.: $\epsilon \times Acc$ computed with GEANT3 and GEANT4 transport codes for deuterons (a) and anti-deuterons (b).

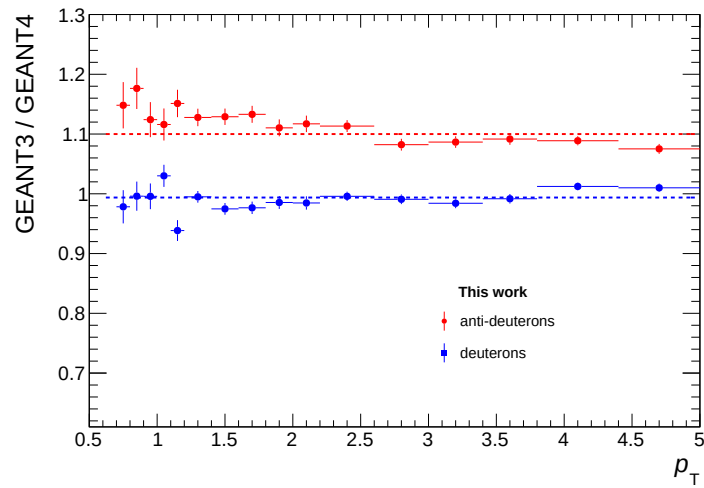


Fig. 4.12.: Ratio between $\epsilon \times Acc$ computed with GEANT3 and GEANT4 transport codes for deuterons. Dashed lines are the mean values.

For the TOF analysis a further correction is necessary, since particles that reach the TOF detector also travelled from the outer TPC wall to the TOF detector, passing through the TRD. In this case, in order to be independent of the MC simulation, a data-driven approach is used to take into account the TPC-TOF matching efficiency. This correction was evaluated for the analysis of the (anti-)deuteron production in the p–Pb data sample collected in 2013 [142]. In that year not all the modules of the TRD were already installed. In this way, it was possible to evaluate the effects of the presence of the TRD by comparing the (anti-)deuteron yield in the region where the TRD modules were present and in those where they were not yet installed. This correction accounts for 1.08 for deuterons and 1.15 for anti-deuterons.

4.7 Fraction of primary deuterons

The fraction of primary deuterons is computed using the methods implemented in the class `TFractionFitter` [143], which is included in the framework ROOT. This class is a tool used for fitting experimental distributions by means of *templates* that are obtained from MC simulations or from real data. A template is a histogram that represents an approximation of a distribution whose analytical expression is unknown. `TFractionFitter` performs fits with the maximum likelihood method. It takes into account the statistical uncertainties of the template by considering the true value of each bin as a fit parameter that can vary in a range defined by its uncertainty. As a result, the shape of the MC distributions can be changed during the fit operations.

Two template functions are used for the fit, one for the primary deuterons and one for the secondary ones. For the DCA_{xy} distribution for primary deuterons a data driven approach is used. Since anti-deuterons are known to be produced only in the primary collision, they can be used to build a realistic DCA_{xy} distribution also for the deuterons. This procedure is justified by the fact that, at the LHC the same production of matter and anti-matter is observed, due to the vanishing baryon-chemical potential. In this way, it is also possible to exclude any bias caused by the injection of deuterons in the MC sample (see Sec. 4.2). The distribution of secondary deuterons is instead taken from the MC sample. However, pp collisions are characterised by a charged particle multiplicity which is much lower than in Pb–Pb collisions and hence also very few secondary deuterons are produced. For this reason, the MC sample with additional injected nuclei are not sufficient to have a template and other 500 million MC events without injected nuclei are used. However, even with all these simulated events, the templates for secondary deuterons for the single multiplicity classes are scarcely populated and hence affected by large statistical uncertainties. However, they are compatible with those obtained from the

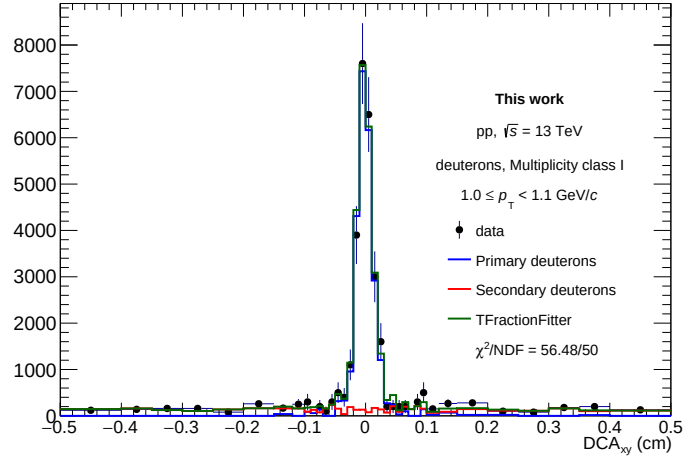


Fig. 4.13.: Fit to DCA_{xy} distribution of deuterons in the MB data sample for $1 \leq p_T < 1.1$ GeV/c. The fit function (green) is the sum of two components, i.e. the template for secondary deuterons (red) and the template for primary deuterons (blue).

MB data sample, which is affected by lower statistical uncertainties. For this reason, for each multiplicity class the templates from secondary deuterons taken from the MB sample are used. On the contrary, the templates for primary deuterons are taken separately for each multiplicity class.

Fig. 4.13 shows a fit to the DCA_{xy} distribution of deuterons in the transverse momentum range $1 \leq p_T < 1.1$ GeV/c for the multiplicity class I. While the distribution for secondary deuterons (red line) is flat, reflecting the absence of a preferential direction in the emission of secondary deuterons, the distribution of primary deuterons (blue line) shows a clear peak around $DCA_{xy} = 0$ cm, as expected from particles produced in the primary vertex. From this distribution, the fraction of primary deuterons $f_{primary}$ is computed as the ratio between the integral of the primary template and the integral of the function that is the sum of the two contributions (green line). Both the integrals are evaluated in the interval $|DCA_{xy}| < 0.12$ cm, because this is the selection applied to the deuteron candidates.

The contamination from secondary deuterons is expected to be higher at low p_T and to decrease and become negligible at high p_T . The fraction of primary deuterons as a function of p_T is shown in Fig. 4.14. The points are obtained from a fit to the DCA_{xy} distribution in different p_T intervals with the method described in the previous paragraph. At high p_T it is not possible to study the DCA_{xy} distribution due to the lack of secondary deuterons and the fraction of primaries must be extrapolated.

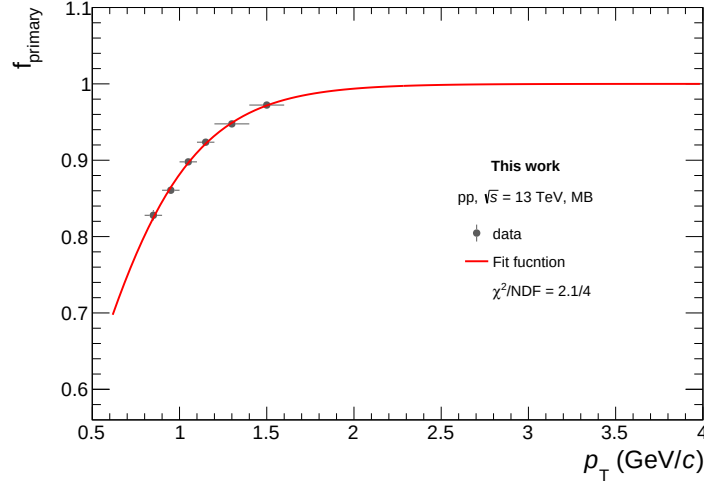


Fig. 4.14.: Fraction of primary deuterons $f_{primary}$ as a function of p_T for the MB analysis.

Data points are fitted with a *ad hoc* function that reproduce the trend and reaches 1 asymptotically at high p_T :

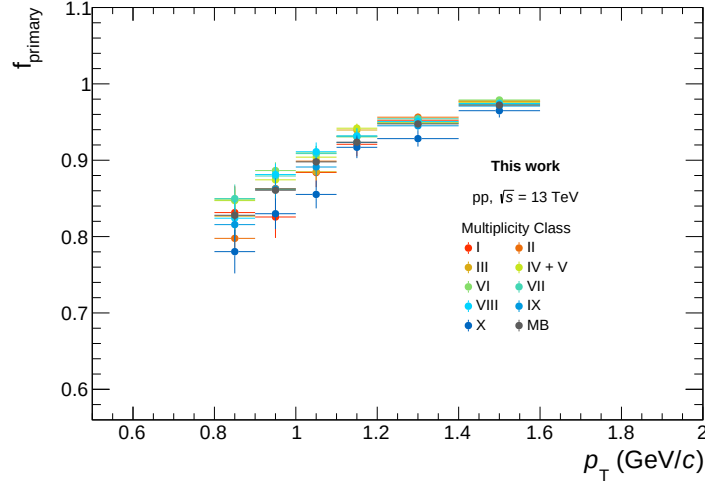
$$f_{primary}(p_T) = \frac{1}{1 - a e^{b p_T}}, \quad (4.10)$$

where a and b are the function parameters. Where the direct evaluation is possible, the point is used, otherwise the fraction of primaries is extrapolated from the fit function. The contribution of secondary deuterons decreases with transverse momentum as expected, becoming negligible at $p_T \approx 2$ GeV/ c . The correction for all the multiplicity classes under study is reported in Fig. 4.15a. It is possible to observe some differences between the corrections computed in the different multiplicity classes: the ratio between $f_{primary}$ evaluated in each multiplicity class and in the MB analysis is shown in Fig. 4.15b. Therefore, for each multiplicity class the specific correction is applied.

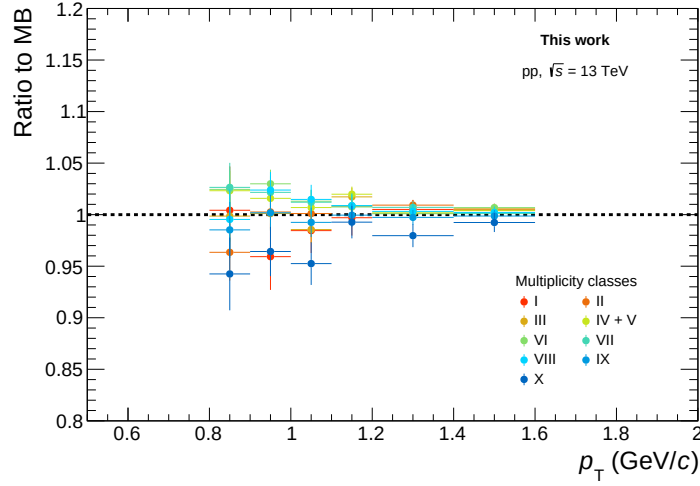
4.8 Event loss and signal loss

As explained in Sec. 4.1, it is necessary to correct the production spectra for both event and signal losses, because of trigger inefficiencies.

The estimation of ϵ_{events} and ϵ_{part} , defined in Eq. 4.3, is computed with a MC simulation, because it is necessary to know whether an event has been erroneously rejected. However, this study cannot be carried out directly on the MC production with injected (anti-)nuclei, because the presence of additional particles would introduce a



(a)



(b)

Fig. 4.15.: Fraction of primary deuterons $f_{primary}$ as a function of p_T for the all the analysed multiplicity classes (a) and ratio with the value obtained in the MB analysis (b).

bias in the estimation of the number of lost (anti-)deuterons. For this reason, the correction is extrapolated by those obtained for pions, kaons and protons from a MC sample without injected nuclei.

The signal loss correction is computed as the ratio between the yields observed in the selected events and that obtained in all the generated events. It is computed for pions (π), kaons (K) and protons (p) in all the analysed multiplicity classes. First, positive and negative particles are considered separately, in order to check whether the signal loss is different for charged-conjugated particles. In Fig. 4.16 this ratio is shown for pions (Fig. 4.16a), kaons (Fig. 4.16b) and protons (Fig. 4.16c). For each particle species and in each multiplicity class the correction factor does not depend on the electric charge of the particle under study. As a consequence, it is not

necessary to separate particles according to the electric charge but charge-conjugated particles can be taken together.

The second step consists in checking whether ϵ_{part} depends on the mass of the particle under study. This check is performed by comparing ϵ_{part} for different species. Fig. 4.17 shows the ratio between ϵ_{part} computed for kaons and that computed for pions (Fig. 4.17a) and the ratio between ϵ_{part} computed for protons and that computed for pions (Fig. 4.17b). In general, these ratios do not show a clear dependence on the mass of the particle, because in some p_T intervals the ratio is above one and in others it is below. Moreover, these ratios differ at most by 1% from the unity. Therefore, the correction is considered mass independent and the mean value of the corrections computed for the different species is applied to (anti-)deuterons.

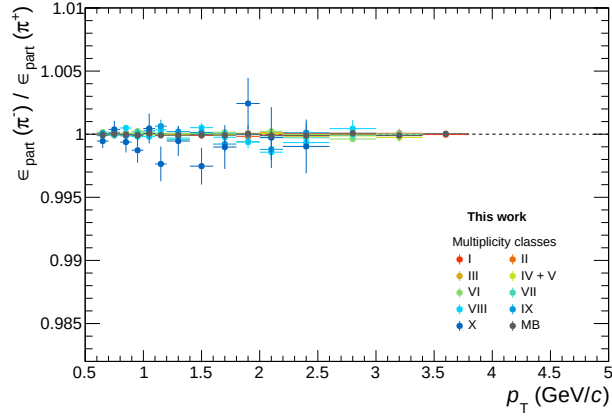
Fig. 4.18 shows the signal loss correction for (anti-)deuterons as a function of transverse momentum in the multiplicity classes used in this work and for the MB analysis. In most of cases, this correction is negligible, but it cannot be neglected for the multiplicity classes VIII, IX and X and for the MB analysis. However, while for multiplicity classes VIII and IX and for the MB this correction is below 3%, for the lowest multiplicity class X the correction is around 14% at low p_T .

Finally, the correction ϵ_{events} to the number of selected events is reported in Tab. 4.8. This correction becomes more relevant going towards low multiplicity. This is expected because, for example, the multiplicity class X has on average less than three reconstructed tracks (see Tab 4.3) and therefore it is more likely that one event is rejected because there are no tracks in the pseudo-rapidity region $|\eta| < 1$. The value for the MB analysis is given by the trigger efficiency and it is lower than the efficiency for INEL >0 , for which a track in the pseudo-rapidity region $|\eta| < 1$ is requested.

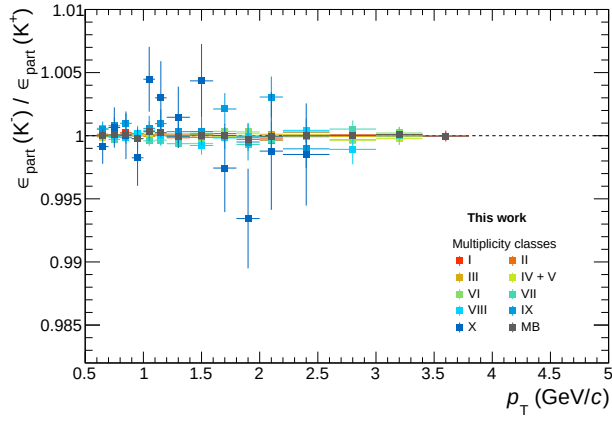
4.9 Systematic uncertainties

The last step of the path from raw yields to production spectra is the evaluation of systematic uncertainties. The relative systematic uncertainties as a function of p_T are shown in Fig. 4.20 for deuterons and in Fig. 4.21 for anti-deuterons. Moreover, the main contribution to systematic uncertainties are reported for Tab. 4.9 for the highest and the lowest multiplicity classes, respectively.

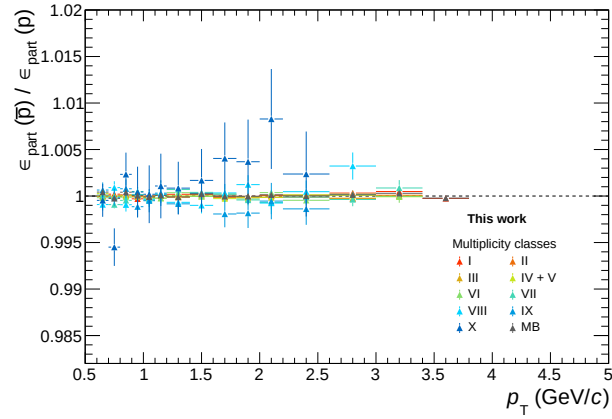
The first contribution to the systematics uncertainties concerns track selections. In particular, by varying the track selection criteria it is possible to take into account



(a) pions



(b) kaons



(c) protons

Fig. 4.16.: Ratio between the ϵ_{part} of charged-conjugated particles as a function of transverse momentum for pions (a), kaons (b) and protons (c).

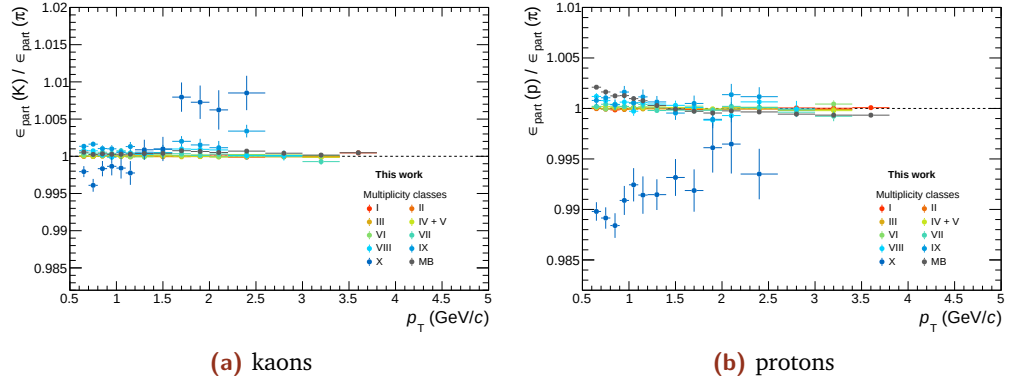


Fig. 4.17.: Ratio between the ϵ_{part} for kaons and that for pions (a) and between ϵ_{part} for protons and that for pions (b).

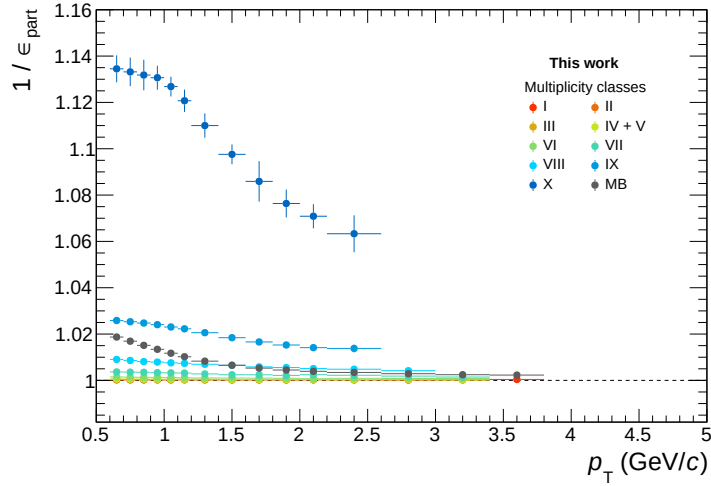


Fig. 4.18.: Signal loss correction for (anti-)deuterons as a function of transverse momentum in the multiplicity classes used in this work and for the MB analysis.

Tab. 4.8.: Event loss correction for all the multiplicity classes and for the MB analysis.

Multiplicity class	ϵ_{events}
I	99.97%
II	99.92%
III	99.94%
IV + V	99.88%
VI	99.68%
VII	99.23%
VIII	98.28%
IX	95.38%
X	78.55%
MB	74.48%

Tab. 4.9.: Summary of the main contributions to the systematic uncertainties for the extreme multiplicity classes I and X. Values in brackets are referred to anti-deuterons. If they are not present, the systematic uncertainty is common for deuterons and anti-deuterons.

Source	d ($\bar{d}$)			
Multiplicity	Class I		Class X	
p_T (GeV/ c)	0.7	3.8	0.7	2.6
Track selection	2% (1%)	2% (3%)	2% (1%)	5% (6%)
Signal extraction	1%	7% (7%)	1%	5% (5%)
Material budget	< 1%	< 1%	< 1%	< 1%
Hadronic interaction	4% (7.5%)	4% (7.5%)	4% (7.5%)	4% (7.5%)
ITS-TPC matching	1%	2.5%	1%	2.5%
Signal Loss	-	-	6%	3%
Total	5% (8%)	9% (11%)	8% (10%)	10% (12%)

the effects that are not properly described in the MC simulation. The selections that are varied as systematic checks are the minimum number of TPC clusters, the TPC pre-selection for the TOF analysis and the selection on the DCA, both in the longitudinal and in the transverse direction. In particular, varying the selections on the DCA is important for checking the correction for secondary deuterons, because the contamination of secondary deuterons is strongly related to this quantity, as discussed in Sec. 4.7. The list of all the applied selections is reported in Tab. 4.10. This source of systematic uncertainty is represented by a red line in Fig. 4.20 and in Fig. 4.21. For each selection variation, the complete analysis is repeated from the start following the same steps described above. By changing the selection criteria, the sample of (anti-)deuterons changes consequently. However, only the statistically significant variations are considered in the systematic uncertainty. In particular, the result obtained by changing the track selections is compared with the standard one and only if the difference is more than 1σ it is taken into account in the systematic uncertainty. This check is generally used in the ALICE collaboration and it is called *Barlow's criterion*, from the paper by Roger Barlow in which it is described [144]. For each source, the systematic uncertainty is evaluated as the standard deviation of the significant variations and the total systematic uncertainty from track selections is the sum in quadrature of these components. The systematic uncertainty from track selection is different from one multiplicity class to another, but it generally increases with transverse momentum. It is interesting to notice that this systematic uncertainty at low p_T is slightly higher for deuterons than for anti-deuterons. This is due to the variation of the selection on the DCA, that is sensitive to the contamination of secondary particles, which is absent in the case of anti-deuterons.

Tab. 4.10.: Summary of the track selections for the evaluation of the systematic uncertainties.

Selection	Default value	Other values
# TPC clusters	70	60, 65, 75, 80
PID (σ_{TPC})	3	3.25, 3.5
DCA _{xy} (mm)	1.2	1.0, 1.4
DCA _z (cm)	1	0.5, 0.75, 1.5, 2

The second source of systematic uncertainty comes from the signal extraction. This systematic uncertainty is evaluated through the direct count of the candidates. This operation is also called *bin counting*. For the TPC analysis bin counting is performed by default in the region $[\mu - 3\sigma, \mu + 3\sigma]$, where μ and σ are the parameters of the fit function used for the signal extraction. The possible miscounting is evaluated by varying the width of the bin-counting interval by 1 or 2 bins and by shifting the interval by 1 or 2 bins. Where background is present, first it is fitted in the

side-bands, i.e. in the complementary region to that in which bin counting is performed. Then, the integral of the background function evaluated in the bin-counting region is subtracted. The related systematic uncertainty is computed as the standard deviation of the different values. Similarly, for the TOF analysis the default interval is $[\mu - 3\sigma, \mu + 5\sigma]$, which is asymmetric to take into account the exponential tail that characterises the signal shape, which is more pronounced than in the TPC analysis. The bin-counting interval is widened and shifted by 1 or 2 bins and the systematic uncertainty is computed as the standard deviation of the different values. In Fig. 4.20 and in Fig. 4.21 it is possible to see that the yield-extraction systematic uncertainty is higher for the TOF analysis ($p_T \geq 1$ GeV/ c) than for the TPC analysis ($p_T < 1$ GeV/ c). This correction is different in each multiplicity class, but in general for the TPC analysis it contributes for less than 1%, while for the TOF analysis it is about 3% where no background is present ($1 \leq p_T < 2.5$ GeV/ c) and, for high multiplicity classes, it increases at higher p_T becoming one of the dominant contributions.

The third source of systematic uncertainty comes from the imperfect reproduction of the detector material budget in MC simulations. This check is performed by computing the $\epsilon \times Acc$ correction using MC simulations in which the material budget is changed from its nominal value. For this study two MC production are used, one with the detector material budget increased by 4.5% and the other with the detector material budget decreased a 4.5%. The value 4.5% corresponds to the uncertainty on the determination of the material budget by measuring photon conversions. $\epsilon \times Acc$ for the MC simulation with varied material budgets are shown in Fig. 4.19. The variation induced by the different material budget is supposed to follow a uniform distribution and therefore the corresponding systematic uncertainty is computed as:

$$\sigma_{Material\ Budget}(p_T) = \frac{\epsilon_{max}(p_T) - \epsilon_{min}(p_T)}{\sqrt{12}}. \quad (4.11)$$

In Fig. 4.20 and in Fig. 4.21 it is possible to see that this source of systematic uncertainty accounts for less than 1%, both for deuterons and anti-deuterons.

The fourth source of systematic uncertainty is given by the incomplete knowledge on the hadronic interaction cross section for (anti-)deuterons, as described in Sec. 4.6.1. This is evaluated from the data driven approach, described in Sec. 4.6.1, in which the yields of (anti-)deuterons obtained in the region where the TRD modules were not installed are compared with the yields obtained in regions where the TRD modules were already installed. The systematic uncertainty related to the hadronic interaction is taken as half of this correction for the whole p_T range. The hadronic interaction systematic uncertainty accounts for 4% for deuterons and 7.5% for anti-deuterons.

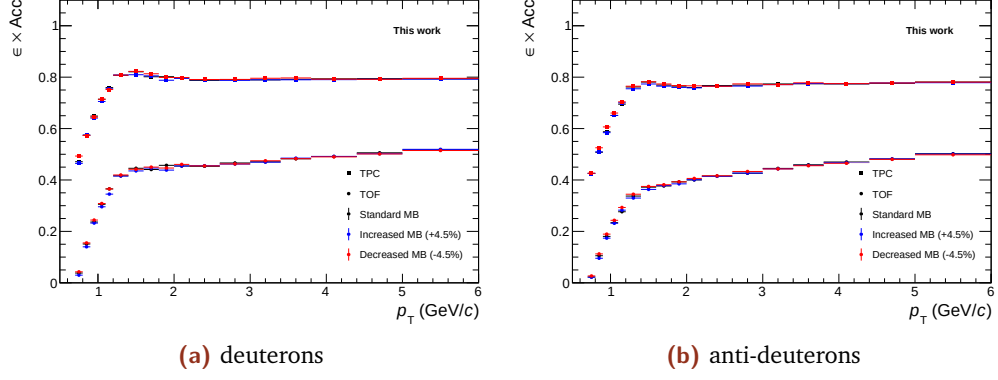


Fig. 4.19.: $\epsilon \times Acc$ computed from MC simulations with nominal, 4.5% increased and 4.5% decreased detector material budget (MB) for deuterons (a) and for anti-deuterons (b).

The fifth source of systematic uncertainty is related to the matching between tracks in the ITS and in the TPC and how it is reproduced in MC. This systematic uncertainty is assumed to be the same for all particle species and it is evaluated centrally within the ALICE Collaboration by comparing the ITS-TPC matching efficiencies in data and MC. This correction depends on p_T : 1% for $p_T < 1$ GeV/c, 1.5% for $1 \leq p_T < 2$ GeV/c and 2% for $p_T > 2$ GeV/c.

Finally, the sixth and last source of systematic uncertainty is given by the signal loss correction, described in Sec. 4.8. It is assumed to be half of the difference between the signal-loss correction and 1. It is strongly dependent on the event multiplicity: it is negligible at high multiplicity (multiplicity classes from I to VII) and contributes up to 6% in the lowest multiplicity class (class X). Where present, it decreases with p_T .

For the computation of the integrated yield dN/dy and the mean transverse momentum $\langle p_T \rangle$ (see Sec. 5.3), it is necessary to estimate the p_T -correlated and the multiplicity-correlated fractions of the systematic uncertainty, as described in Sec. 4.9.1 and in Sec. 4.9.2, respectively.

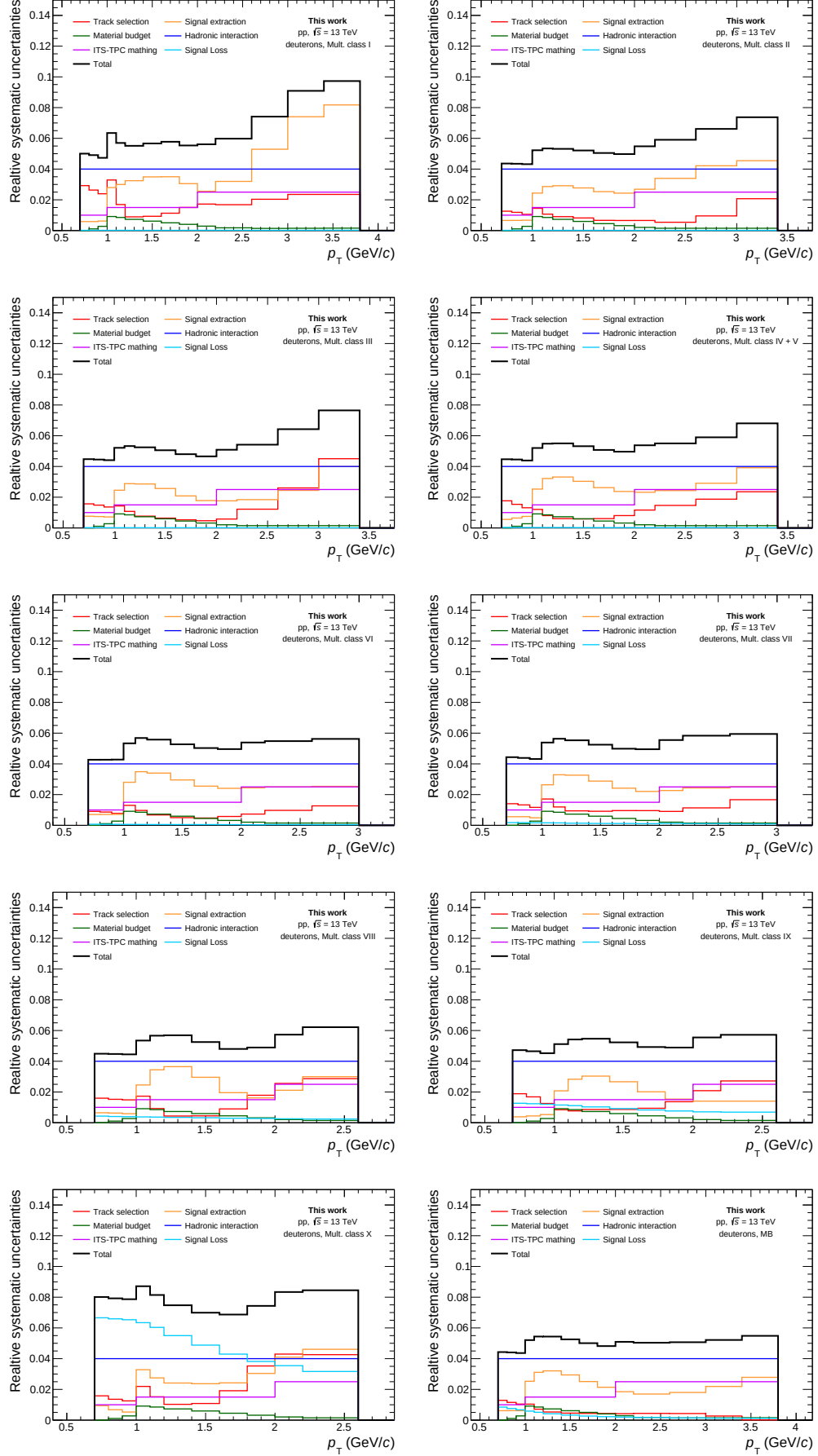


Fig. 4.20.: Relative systematic uncertainties as a function of p_T for deuterons in all the analysed multiplicity classes and in MB analysis.

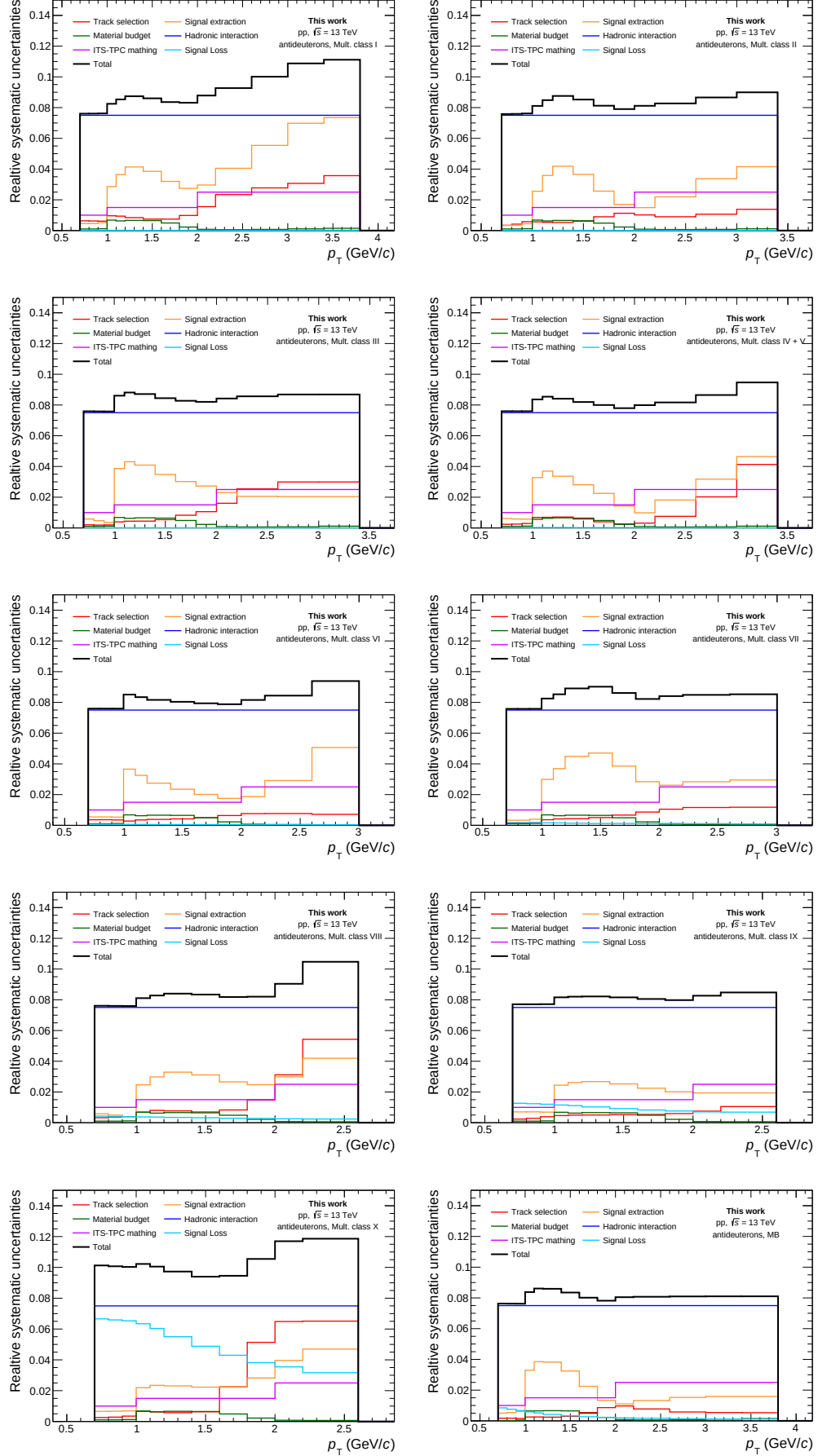


Fig. 4.21.: Relative systematic uncertainties as a function of p_T for anti-deuterons in all the analysed multiplicity classes and in MB analysis.

4.9.1 Determination of the transverse-momentum-correlated systematic uncertainty

The list of the fraction of p_T -correlation for each source of systematic uncertainty is reported in Tab. 4.11. Some sources are assumed to be fully p_T -correlated. For example, the uncertainty on the material budget is assumed to be fully correlated because, if the material budget is varied, it affects all the p_T intervals in the same direction. Also the uncertainty related to hadronic interaction is assumed to be fully correlated, because different definitions of the cross section affect in the same direction the yields in all the p_T intervals. For the same reason, the ITS-TPC matching is assumed as fully p_T -correlated. On the contrary, signal loss systematic uncertainty is assumed to be fully uncorrelated, because this uncertainty has been evaluated individually for each p_T value and it is different for each multiplicity class.

For the other sources, the fraction of p_T -correlation is expressed as the correlation coefficient of the scatter plots with the spectra obtained with the different selections. The scatter plots are reported in Fig. 4.22. Considering the MB analysis, for each variation of the selection criterium the ratio between the varied yield S_{var} and the reference one S_{ref} is computed for each p_T bin and plotted at the corresponding value of the applied selection. In this way, for each value of the selection variable a distribution of S_{var}/S_{ref} is obtained. Finally, from the scatter plot the correlation coefficient is computed and this value is taken as the p_T -correlated fraction of the systematic uncertainty.

The sources that show the highest fraction of p_T -correlation are the selection on the DCA, which is included in the track-selection uncertainty, and the broadening of the range in which the bin counting is performed, included in the signal-extraction uncertainty. However, except for few cases, these contributions are less significant than the hadronic interaction uncertainty, which is fully p_T -correlated. Therefore, it is possible to state that most of the total systematic uncertainty is p_T -correlated.

4.9.2 Determination of the multiplicity-correlated systematic uncertainty

Concerning the multiplicity-correlated systematic uncertainty, some sources are assumed to be fully multiplicity-correlated, because they are the same in all the multiplicity classes. In particular, the material budget, the hadronic interaction and the ITS-TPC matching systematic uncertainties are fully correlated. The signal loss uncertainty is fully uncorrelated, because it is evaluated independently for each multiplicity class.

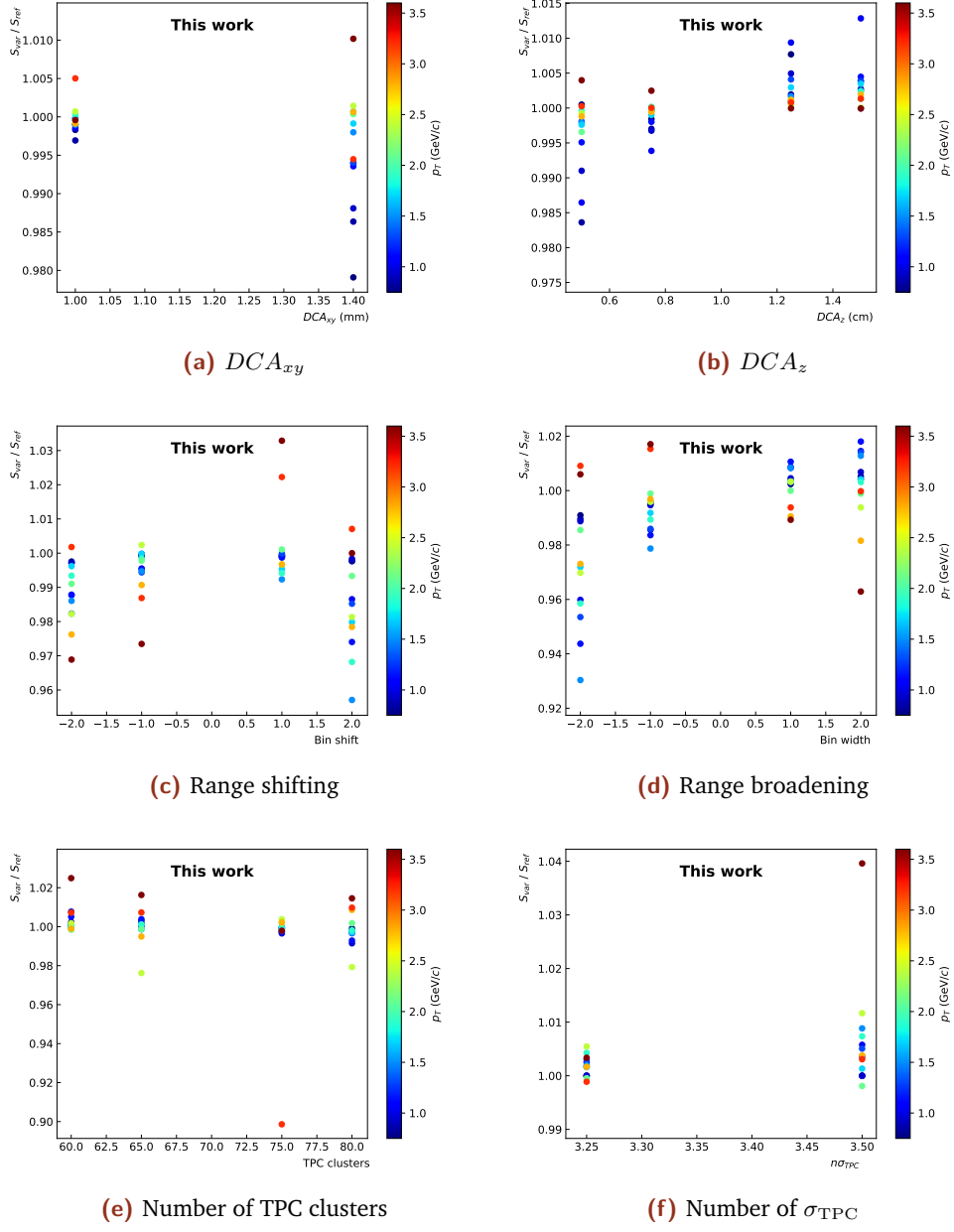


Fig. 4.22.: Scatter plot of the ratio between the varied yield S_{var} and the reference one S_{ref} for different selection criteria.

Tab. 4.11.: p_T -correlated fraction for each source of systematic uncertainty.

Source	p_T -correlated fraction (%)
Material budget	100
Hadronic interaction	100
ITS-TPC matching	100
Signal loss	0
DCA_{xy}	34
DCA_z	58
TPC clusters	18
PID	33
Width bin counting	57
Shift bin counting	7

For the other sources, the fraction is evaluated through the R ratio, defined as:

$$R_{var}^i = \frac{\frac{S_{var}^i}{S^i}}{\frac{S_{var}^{MB}}{S^{MB}}} = \frac{1 + \frac{\Delta S_{var}^i}{S^i}}{1 + \frac{\Delta S_{var}^{MB}}{S^{MB}}} \quad (4.12)$$

where R_{var}^i is the R coefficient for the multiplicity class i corresponding to the variation var under study. S^i and S^{MB} are the nominal spectra for the multiplicity class i and the MB sample, respectively. S_{var}^i and S_{var}^{MB} are the varied spectra, corresponding to the variation var under study, for the multiplicity class i and the MB sample, respectively. Basically this ratio compares the variation of the spectra in a specific multiplicity class with that obtained in the MB sample. In the third term of the equation, the explicit terms in numerator and denominator are the systematic uncertainties corresponding to the variation var .

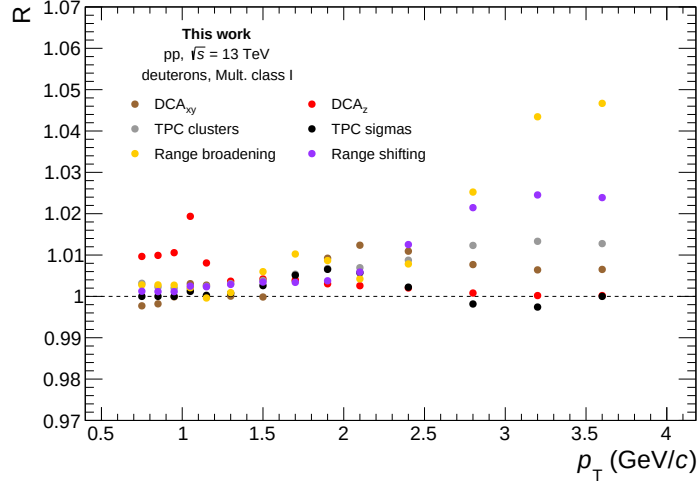
If the uncertainty is fully correlated, then $R_{var}^i = 1$. For this reason, the fraction of uncorrelated uncertainty is evaluated as:

$$f_{uncorr} = |1 - R_{var}^i| \quad (4.13)$$

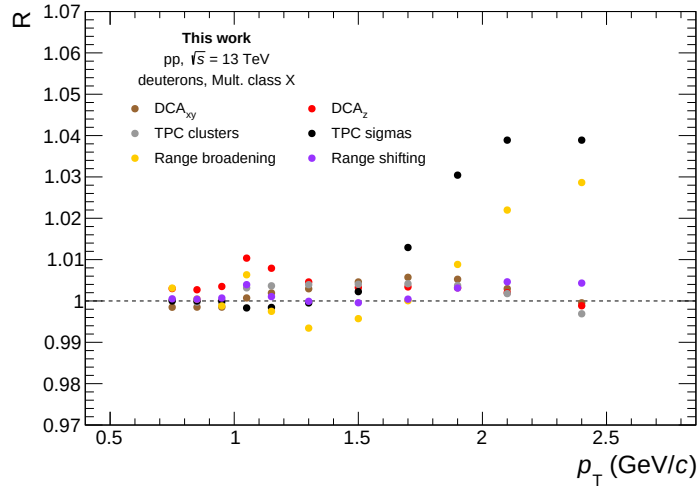
The correlated fraction is complementary:

$$f_{corr} = 1 - |1 - R_{var}^i| \quad (4.14)$$

R values for the two extreme multiplicity classes are shown in Fig. 4.23. One can see that in general these values are very close to the unity, hence the uncertainties are almost fully multiplicity correlated.



(a) Multiplicity class I



(b) Multiplicity class X

Fig. 4.23.: R value for different sources of systematic uncertainty for the two extreme multiplicity classes.

(Anti-)deuteron production spectra

This chapter is dedicated to the description of the results obtained for the production spectra of deuterons and anti-deuterons. The two p_T spectra will be compared and the p_T -integrated yields dN/dy will be computed, together with the mean transverse momentum $\langle p_T \rangle$, both in multiplicity classes and for the minimum bias data sample.

5.1 Production spectra

In Fig. 5.1 the p_T spectra of deuterons (Fig. 5.1a) and anti-deuterons (Fig. 5.1b) are shown. They are shown for the analysed multiplicity classes and for the minimum bias data sample. The p_T spectra are measured in nine multiplicity classes, reported in Tab. 4.3. This is the first measurement of (anti-)deuteron production in pp collisions by the ALICE Collaboration with so many multiplicity classes. The fine multiplicity binning allows for a more differential insight on the multiplicity dependence of the (anti-)deuteron production. Moreover, the measurement is performed up to the same charged particle multiplicity ($dN_{ch}/d\eta \sim 26$) as measured in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142]. The production spectra are obtained from the respective raw spectra by applying all the corrections described in chapter 4, according to Eq. 4.2. For visibility, all the spectra are scaled by powers of 2.

As explained in Sec. 4.5, the signal extraction is performed with the TPC or with the TOF according to the transverse momentum. The final production spectra are obtained by putting together the spectra obtained with TPC and those obtained with TOF: the TPC is used for $p_T < 1$ GeV/ c and the TOF is used for $p_T \geq 1$ GeV/ c . The watershed is $p_T = 1$ GeV/ c because for $p_T > 1$ GeV/ c the background for the TPC is not negligible (see Fig. 4.4b) and $p_T < 1$ GeV/ c the TOF efficiency is quite low (see Fig. 4.9). In the interval $0.9 \leq p_T < 1.2$ GeV/ c , the signal can be extracted with both the methods. Fig. 5.2 shows the overlap of the TPC and the TOF spectra in the aforementioned p_T interval for deuterons (Fig. 5.2a) and anti-deuterons (Fig. 5.2b). The TPC and the TOF spectra are compared for deuterons (Fig. 5.2c) and for anti-deuterons (Fig. 5.2d). In Fig. 5.2c and in Fig. 5.2d, the variable on the y-axis is the difference between the corrected yield obtained with the TPC (N_{TPC})

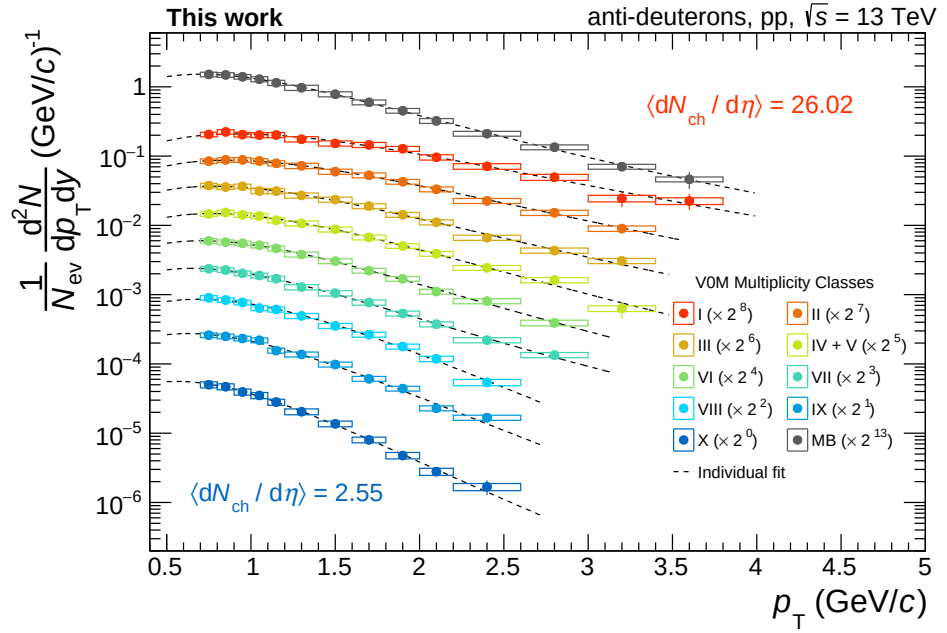
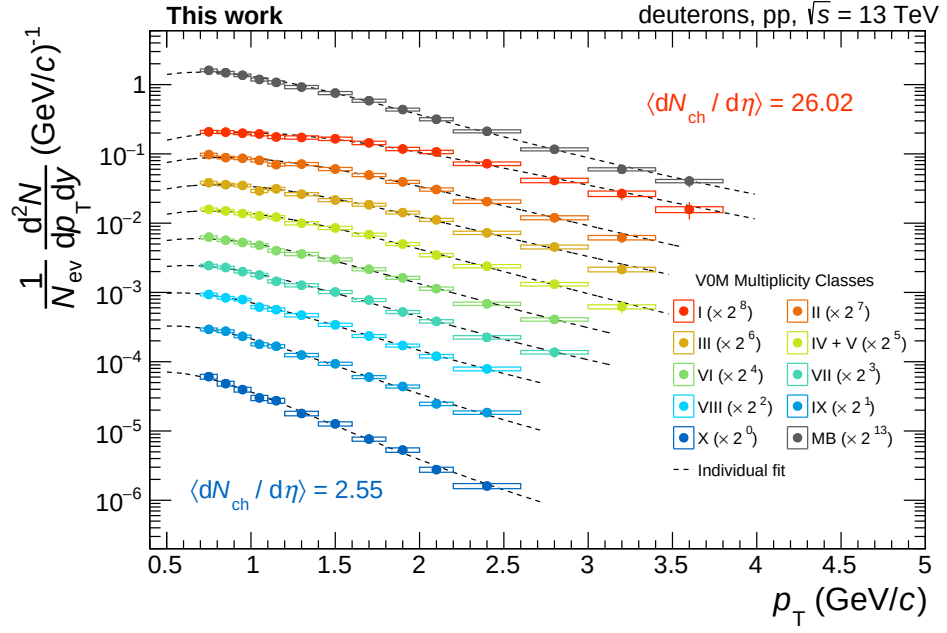


Fig. 5.1.: Transverse momentum spectra for deuterons (a) and for anti-deuterons (b) in the multiplicity classes under study and for the MB analysis. Spectra are scaled for visibility and are fitted with a Lévy-Tsallis function.

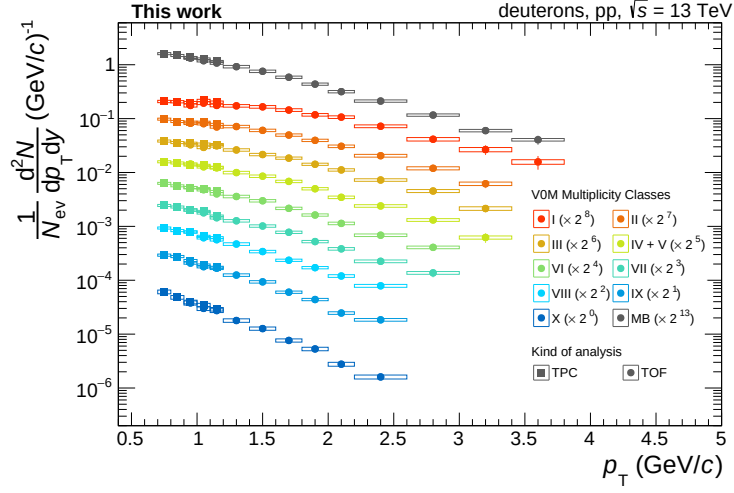
and the yield obtained with the TOF (N_{TOF}), divided by σ , defined as the sum in quadrature of statistical and systematic uncertainties of N_{TPC} and N_{TOF} . Except for some small fluctuations in the deuteron spectra, the corrected yields obtained with the TPC analysis and with the TOF analysis are always in agreement within 1σ . The matching is more precise for anti-deuterons for two reasons. First, the uncertainties on the anti-deuteron yields are larger than those on the deuteron yield. Second, the correction for secondary particles is applied only to deuterons (see Sec. 4.7) and it is an additional possible source of disagreement.

Fig. 5.1 shows that going from low multiplicity (blue points) to high multiplicity (red points) the spectra become less and less steep. This behaviour was already observed in Pb-Pb collision at $\sqrt{s_{NN}} = 2.76$ TeV [80], in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] and in pp collisions at $\sqrt{s} = 7$ TeV [85]. This trend is more evident in Pb-Pb collisions than in small systems and it is interpreted as an indication of the presence of radial flow, that modifies the shapes of the spectra according to the pressure gradients present in the expanding source. Even if less pronounced, this effect can also be seen in this analysis and can be interpreted as a hint of radial flow in pp collisions. This trend is more clear in this analysis than in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] and in pp collisions at $\sqrt{s} = 7$ TeV [85] thanks to the finer multiplicity binning and to the wider p_T range.

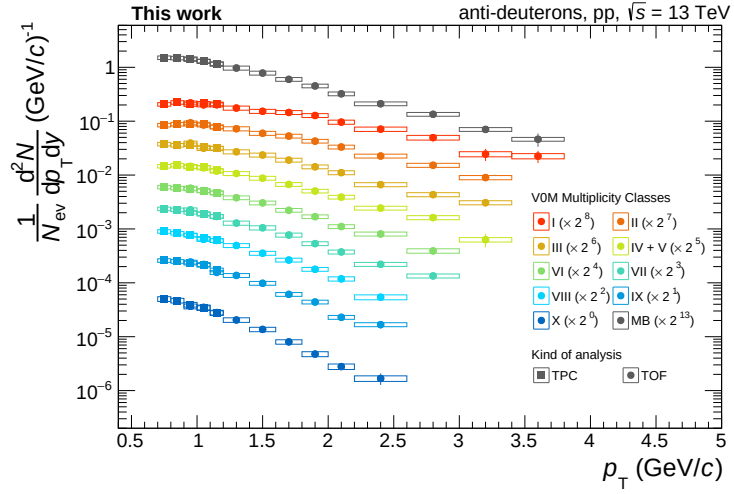
The p_T spectra are fitted with a Lévy-Tsallis function [145] to extrapolate yields outside the measured p_T region for the computation of the p_T -integrated yields and of the mean transverse momentum $\langle p_T \rangle$. The Lévy-Tsallis function is defined as:

$$\frac{d^2N}{dy dp_T} = \frac{dN}{dy} \frac{p_T (n-1)(n-2)}{nC[nC + m_d(n-2)]} \left(1 + \frac{m_T - m_d}{nC}\right)^{-n}, \quad (5.1)$$

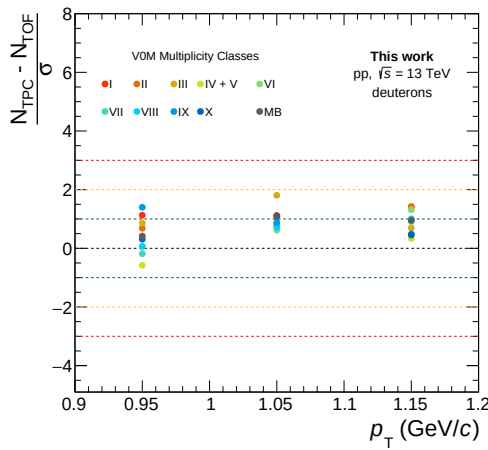
where m_d is the deuteron mass, $m_T = \sqrt{m_d^2 + p_T^2}$ is the transverse mass, n , C and dN/dy are free fit parameters. The parameters of the fit function for the deuteron and the anti-deuteron spectra are reported in Tab. 5.1. As it will be discussed in Sec. 5.3, the fit function is used to extrapolate the yields in the unmeasured region at low and high p_T . This step is necessary for the computation of the p_T -integrated yield dN/dy and of the mean transverse momentum $\langle p_T \rangle$.



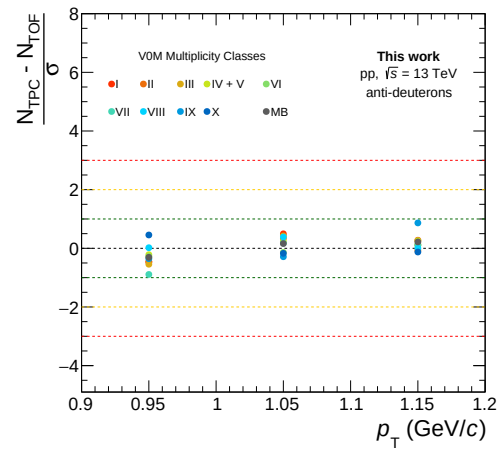
(a)



(b)



(c)



(d)

Fig. 5.2.: Superposition of p_T spectra obtained with TPC (squares) and TOF (circles) for deuterons (a) and anti-deuterons (b) for the different multiplicity classes and for the MB analysis. The yields in the p_T intervals which are common for the TPC and the TOF analyses are compared in (c) and (d).

Tab. 5.1.: Parameters of the Lévy-Tsallis function used to fit the deuteron spectra (a) and the anti-deuteron spectra (b).

(a) deuterons				
Mult. class	n	C (GeV)	$dN/dy (\times 10^{-4})$	χ^2 / NDF
I	10 ± 2	0.40 ± 0.04	15.4 ± 0.4	0.63
II	8 ± 2	0.31 ± 0.03	11.7 ± 0.3	1.11
III	6 ± 2	0.28 ± 0.02	9.1 ± 0.2	0.97
IV+V	7 ± 2	0.25 ± 0.02	7.1 ± 0.2	0.61
VI	6 ± 1	0.21 ± 0.02	5.3 ± 0.2	0.76
VII	4.7 ± 0.8	0.17 ± 0.02	4.1 ± 0.2	0.85
VIII	4.4 ± 0.8	0.15 ± 0.02	3.1 ± 0.1	0.35
IX	4.7 ± 0.8	0.13 ± 0.02	1.85 ± 0.09	1.08
X	5 ± 1	0.11 ± 0.02	0.69 ± 0.06	0.55
MB	5.7 ± 0.6	0.22 ± 0.01	2.67 ± 0.06	0.79

(b) anti-deuterons				
Mult. class	n	C (GeV)	$dN/dy (\times 10^{-4})$	χ^2 / NDF
I	7 ± 3	0.37 ± 0.05	16.1 ± 0.8	0.32
II	7 ± 3	0.32 ± 0.04	12.2 ± 0.5	0.08
III	6 ± 2	0.27 ± 0.03	9.4 ± 0.4	0.23
IV+V	8 ± 3	0.27 ± 0.03	7.1 ± 0.3	0.21
VI	7 ± 2	0.23 ± 0.03	5.3 ± 0.2	0.17
VII	6 ± 2	0.19 ± 0.03	4.0 ± 0.2	0.15
VIII	17 ± 13	0.23 ± 0.03	2.7 ± 0.1	0.34
IX	10 ± 5	0.19 ± 0.03	1.63 ± 0.08	0.76
X	10 ± 5	0.15 ± 0.02	0.59 ± 0.05	0.11
MB	6 ± 1	0.22 ± 0.02	2.73 ± 0.09	0.11

5.2 Ratio between anti-deuteron and deuteron spectra

The most direct way to compare deuteron and anti-deuteron spectra is to compute the ratio of the two. From previous measurements in different collision systems and at different energies, the production spectra of anti-nuclei and nuclei have always been found to be compatible [80, 85, 142], as well as all the spectra of charged-conjugated particles.

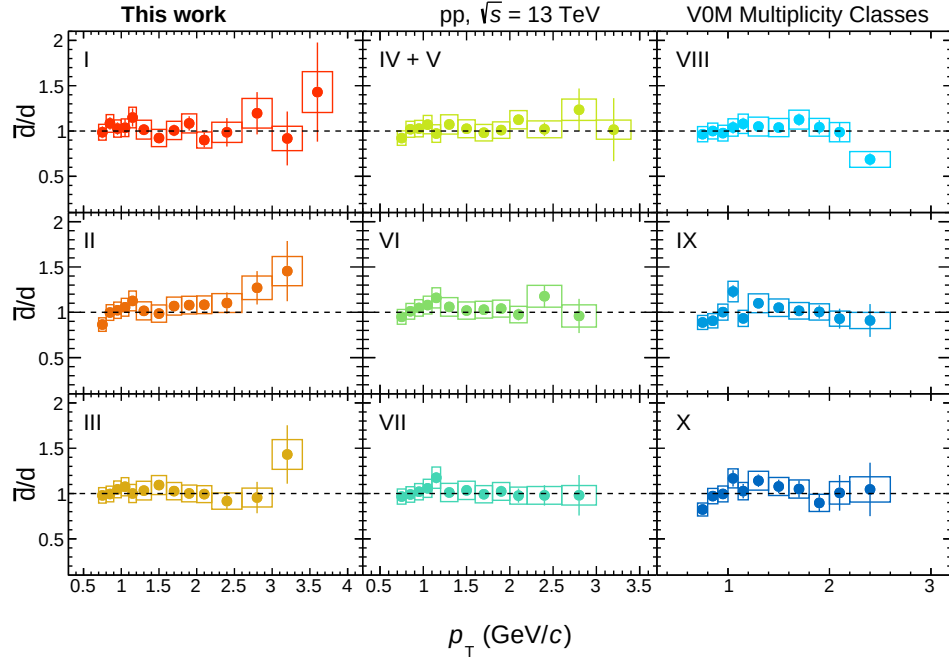
According to both the statistical hadronisation model and the coalescence model, the ratio between the yields of anti-nuclei and nuclei with mass number A is related to the ratio between anti-proton and proton yields through the relation:

$$\frac{N_{\bar{A}}}{N_A} = \left(\frac{N_{\bar{p}}}{N_p} \right)^A, \quad (5.2)$$

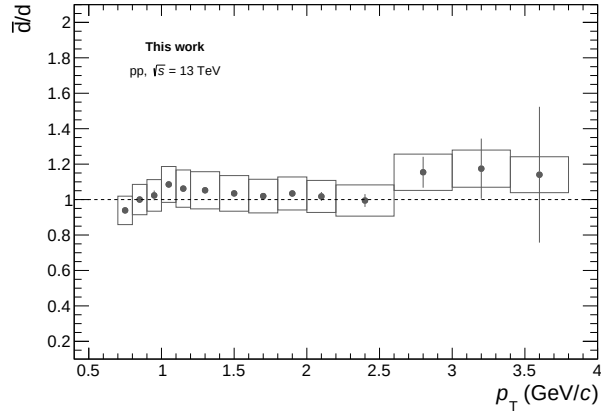
where $N_{\bar{A}}$ and N_A are the number of anti-nuclei and nuclei with mass number A respectively and $N_{\bar{p}}$ and N_p are the number of anti-protons and protons respectively. Since at the LHC the production spectra of protons and anti-protons are compatible, the ratio between the yields of anti-nuclei and nuclei is expected to be compatible with unity too.

Fig. 5.3 shows the ratio between the anti-deuteron and the deuteron spectra for the analysed multiplicity classes (Fig. 5.3a) and for the MB analysis (Fig. 5.3b). The ratios as a function of transverse momentum has also been fitted with a zeroth order polynomial and the results are reported in Tab. 5.2, together with the value of the relative reduced χ^2 . In all the cases, the ratio between the anti-deuteron and the deuteron spectra is compatible with unity, as foreseen from both the statistical hadronisation model and the coalescence model.

In conclusion, it has been verified that in pp collisions at $\sqrt{s} = 13$ TeV there is the same production of deuterons and antideuterons, as already observed in past analysis.



(a) Multiplicity classes



(b) Minimum Bias analysis

Fig. 5.3.: Ratio between the production spectra of anti-deuterons and deuterons for all the analysed multiplicity classes (a) and for the Minimum Bias analysis (b).

Tab. 5.2.: Ratios between anti-deuterons and deuterons spectra for all the multiplicity classes studied in this analysis as a function of p_T . The values are obtained with a fit with a zeroth order polynomial.

Multiplicity class	$\bar{d} / d$	χ^2/NDF
I	1.01 ± 0.04	4.4 / 13
II	1.02 ± 0.03	8.4 / 12
III	1.01 ± 0.03	3.3 / 12
IV + V	1.01 ± 0.03	3.7 / 12
VI	1.03 ± 0.03	3.4 / 11
VII	1.01 ± 0.03	2.2 / 11
VIII	0.99 ± 0.03	10.3 / 10
IX	0.98 ± 0.03	7.7 / 10
X	0.99 ± 0.04	8.8 / 10
MB	1.03 ± 0.03	3.4 / 13

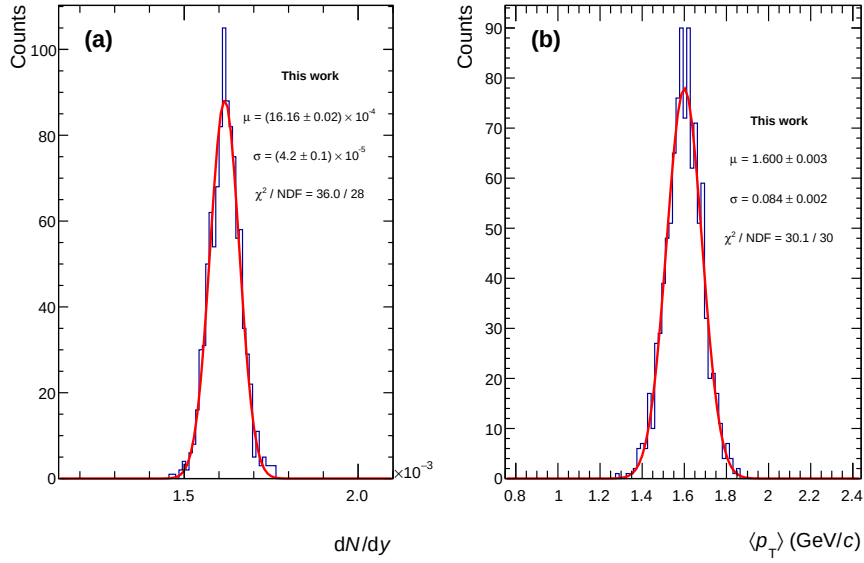


Fig. 5.4.: Distributions of dN/dy (a) and $\langle p_T \rangle$ (b) as obtained in the procedure for the evaluation of the statistical uncertainties (see text for details) for anti-deuterons in the multiplicity class I.

5.3 Integrated yield and mean transverse momentum

The measurement of the p_T -integrated yield dN/dy and the mean transverse momentum $\langle p_T \rangle$ as a function of multiplicity is important to study the dependence of production of (anti-)deuterons and of their kinematics on the system size. For each multiplicity class under study, dN/dy and $\langle p_T \rangle$ are computed relying both on the spectra themselves and on the fit function. dN_{ch}/dy is computed as the sum of two parts. The first is the integral of the measured p_T spectra. The second is the integral of the fit function in the unmeasured region at low and high p_T . Similarly, $\langle p_T \rangle$ is evaluated using both the measured yields and the extrapolated values at low and high p_T .

The statistical uncertainties on dN/dy and $\langle p_T \rangle$ are evaluated by shifting data points by a random value from a Gaussian distribution, centred in 0 and with a standard deviation that is equal to the statistical uncertainty. dN/dy and $\langle p_T \rangle$ are evaluated for the new modified distribution using the standard analysis method and this process of shifting is repeated 1000 times. In this way, two distributions, one for dN/dy and one for $\langle p_T \rangle$, are obtained and they are fitted with a Gaussian, as shown in Fig. 5.4. The σ parameter of the Gaussian is taken as statistical uncertainty.

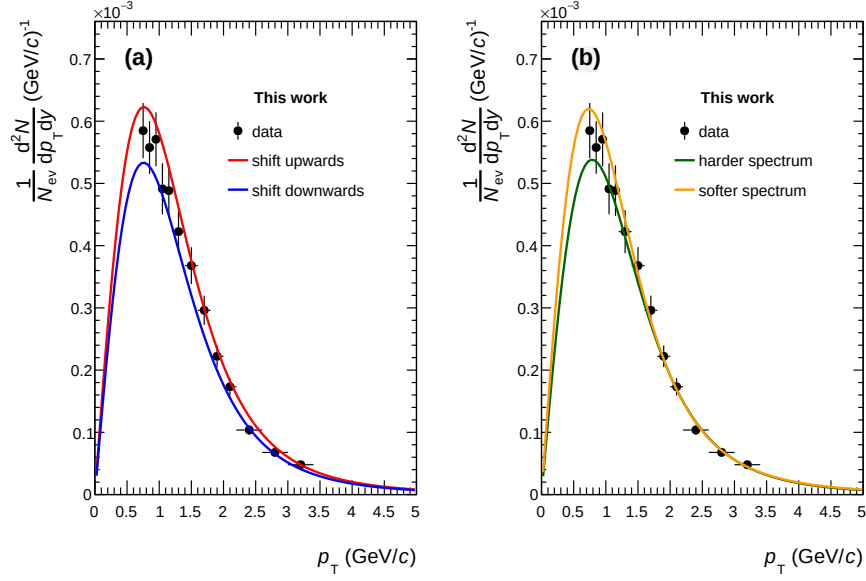


Fig. 5.5.: Anti-deuteron production spectra with p_T -correlated systematic uncertainties in the multiplicity class III. The red (blue) line in panel (a) is a Lévy-Tsallis function used to fit the p_T spectra shifted upwards (downwards) while the green (orange) line in panel (b) is a Lévy-Tsallis function used to fit the p_T spectra made harder (softer) (see the text for details).

The systematic uncertainties are propagated from the p_T spectra (see Sec. 4) and their multiplicity-correlated part is separated from the multiplicity-uncorrelated one (see Sec. 4.9.2). This is the first analysis about the production of nuclei by the ALICE Collaboration in which these two parts are separated. In order to quantify these two kinds of uncertainties one has to first separate the p_T -correlated uncertainty of the spectra from the p_T -uncorrelated one (see Sec. 4.9.1). Then, for the p_T -correlated systematic uncertainty of the spectra, the multiplicity-correlated part is computed. The p_T -correlated part of the systematic uncertainties on the spectra is the only one that can have a relevant component of multiplicity-correlated uncertainty because of the integration over p_T .

For the p_T -uncorrelated part of the systematic uncertainty on dN/dy and on $\langle p_T \rangle$, the same procedure described above for the statistical uncertainty is used. For the p_T -correlated part a different approach is used. For the p_T -correlated systematic uncertainty on dN/dy , the spectra are coherently shifted upwards/downwards for a value corresponding to their p_T -correlated systematic uncertainty. Then, a new value of dN/dy is computed for the modified spectrum using the same procedure used for computing the reference value. The difference between the dN/dy obtained with the modified spectrum and the reference one is taken as systematic uncertainty. An example of the p_T -spectra shifted upwards/downwards is shown in Fig. 5.5a.

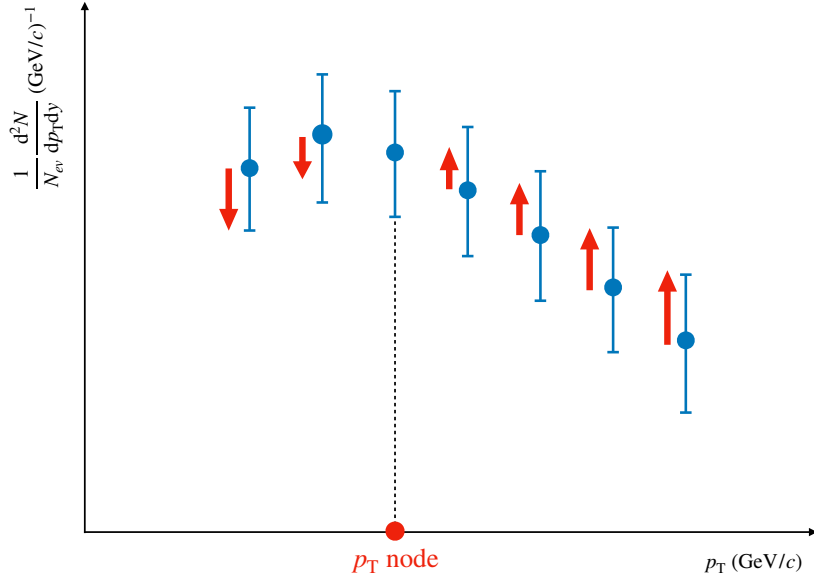


Fig. 5.6.: Scheme of the hardening of a p_T spectrum. The larger is the distance from the p_T node, the larger is the vertical shift (see the text for details).

Similarly, the systematic uncertainty for $\langle p_T \rangle$ is obtained by making the spectra harder/softer according to their systematic uncertainties and following the same procedure used for dN/dy . A scheme of the hardening process is shown in Fig. 5.6. Harder (softer) spectra are obtained as follows: first, a p_T value inside the measured range (defined as *node*) is fixed. The first node chosen corresponds to p_T value of the first point. At the second iteration procedure the p_T node is moved to the adjacent point on the right. After the p_T node has been fixed, all the points with a p_T lower than the node are shifted downwards (upwards) for a fraction of their p_T -correlated systematic uncertainty. The amount of the shift is proportional the p_T -distance between the considered point and the node. On the contrary, all the points with a p_T higher than the node are shifted upwards (downwards), proportionally to their p_T -distance from the node. For the modified spectrum, a $\langle p_T \rangle$ is computed as for the reference value. This procedure is repeated varying the position of the node within the measured p_T range. The difference between the maximum (minimum) $\langle p_T \rangle$ obtained with a harder (softer) spectrum and the reference one is taken as p_T -correlated systematic uncertainty. An example of harder/softer p_T spectra is shown in Fig. 5.5b.

For the multiplicity-correlated systematic uncertainty on dN/dy and $\langle p_T \rangle$, the same procedure used for the p_T -correlated systematic uncertainty is applied, but using the multiplicity-correlated systematic uncertainties instead of the p_T -correlated ones. Note that the multiplicity-correlated systematic uncertainty is contained in the p_T -correlated one.

The last source of systematic uncertainty is determined by the use of different fit functions for the extrapolation in the unmeasured p_T region. Two additional functions are used to this purpose.

The first one is a Boltzmann function, defined as:

$$\frac{d^2N}{dy dp_T} = A \cdot p_T \cdot m_T \cdot e^{-\frac{m_T}{T}}, \quad (5.3)$$

where A and T are fit parameters.

The second one is a m_T -exponential, defined as:

$$\frac{d^2N}{dy dp_T} = A \cdot e^{-\frac{m_T}{T}}, \quad (5.4)$$

where A and T are fit parameters.

For each alternative fit function, dN/dy and $\langle p_T \rangle$ are computed and half of the difference between the maximum and the minimum values is taken as systematic uncertainty from extrapolation. The systematic uncertainty from extrapolation is considered multiplicity uncorrelated, because computed independently for each spectrum, and summed in quadrature to the other multiplicity-uncorrelated systematic uncertainties.

All the measured values of dN/dy and $\langle p_T \rangle$ for deuterons and anti-deuterons are reported in Tab. 5.3 together with the statistical uncertainty, the multiplicity-uncorrelated and the multiplicity correlated systematic uncertainties. It is possible to see how both dN/dy and $\langle p_T \rangle$ decrease going from hi multiplicity (class I) to low multiplicity (class X).

5.4 Comparison with measurements in different collision systems and at different energies

In this section, the measurements reported in Tab. 5.3 are compared with those obtained in previous measurements in different collision systems and at different energies. It is sufficient to compare the measurements for deuterons because the ratio between the deuteron and the anti-deuteron spectra is compatible with 1.

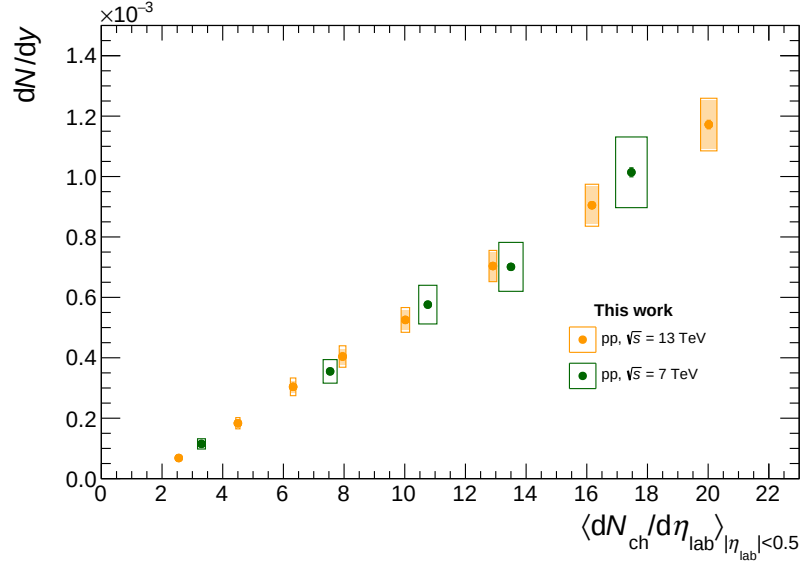
In Fig. 5.7a the dN/dy of deuterons measured in pp collisions at $\sqrt{s} = 13$ TeV (orange points) are compared with those obtained in pp collisions at $\sqrt{s} = 7$ TeV [85] (orange points). As anticipated, thanks to the large data sample, it has been possible

to perform a study of deuteron production that is more granular in multiplicity. Moreover, the systematic uncertainties have also been reduced with respect to the analysis in pp collisions at $\sqrt{s} = 7$ TeV. In addition, in this work the multiplicity correlated part of the systematic uncertainties is put in evidence (shaded area) and it is possible to see that a big part of the total uncertainty is multiplicity correlated. It is interesting to observe that the dN/dy measured at the two energies are in very good agreement. For this reason, it is possible to conclude that the number of (anti-)deuterons produced in a pp collision does not depend on the energy of the collision, but only on the event multiplicity.

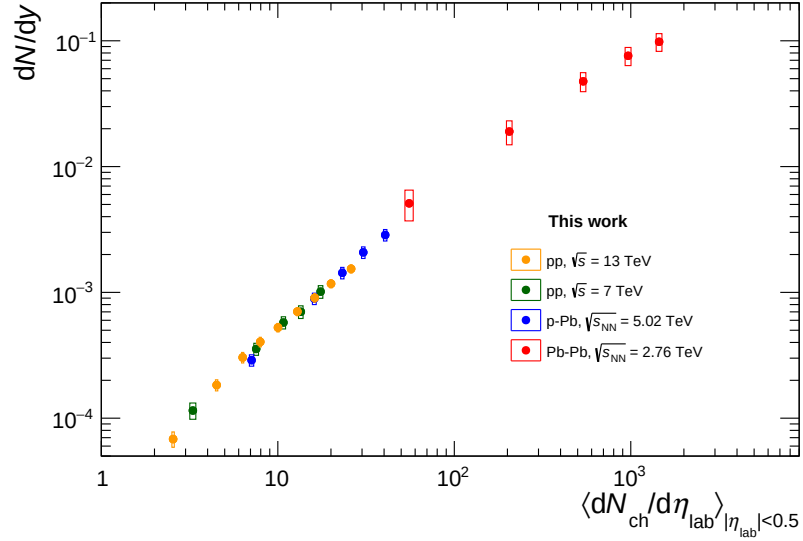
Furthermore, in Fig. 5.7b the deuteron dN/dy is compared across different collision systems. The results are compared with the deuteron dN/dy in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] (blue points) and with the deuteron dN/dy in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80] (red points). dN/dy for (anti-)deuterons evolves smoothly as a function of multiplicity, regardless of the collisions system in which the particles are produced. In other words, the (anti-)deuteron yield can be described as a universal function of charged particle multiplicity, hence of system size.

In Fig. 5.8a, the deuteron $\langle p_T \rangle$ in pp collisions at $\sqrt{s} = 13$ TeV is compared with the $\langle p_T \rangle$ measured in pp collisions at $\sqrt{s} = 7$ TeV [85]. As for dN/dy , also for $\langle p_T \rangle$ the measurements obtained at $\sqrt{s} = 13$ TeV are more precise than those obtained at $\sqrt{s} = 7$ TeV, mainly for a different approach in the treatment of the systematic uncertainty related to signal extraction, that in pp collisions at $\sqrt{s} = 7$ TeV accounted for about 25%, to be compared with the 5% of this analysis. The multiplicity correlated part of the systematic uncertainty is put in evidence for the results of this work (shaded area). In this case, the contributions from the multiplicity-uncorrelated part of the systematic uncertainties is larger than in dN/dy . The reason is that for $\langle p_T \rangle$ the systematic uncertainties coming from the extrapolation with different fit functions is larger. The observation that the impact of a different fit function is larger for $\langle p_T \rangle$ than for dN/dy means that the relative change of the mean value of the function is larger than the relative change of its integral.

In Fig. 5.8b, the deuteron $\langle p_T \rangle$ is shown for pp collisions at $\sqrt{s} = 13$ TeV (orange points), pp collisions at $\sqrt{s} = 7$ TeV (green points), p-Pb collisions at $\sqrt{s_{NN}} = 5$ TeV [142] (blue points) and Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80] (red points). Mean transverse momentum evolves smoothly as a function of multiplicity passing from pp collisions to p-Pb collisions. In particular, for the two colliding systems the measured values of $\langle p_T \rangle$ are compatible. The mean transverse momentum evolves smoothly as a function of multiplicity also in Pb-Pb collisions, but there is a discontinuity going from small systems, i.e. pp and p-Pb collisions, to Pb-Pb collisions. This behaviour was already observed for other light-flavoured hadrons, as



(a)

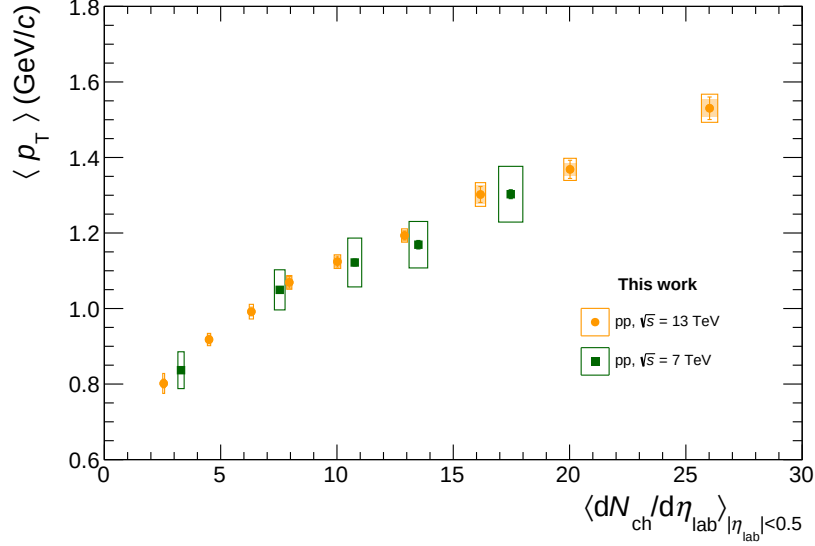


(b)

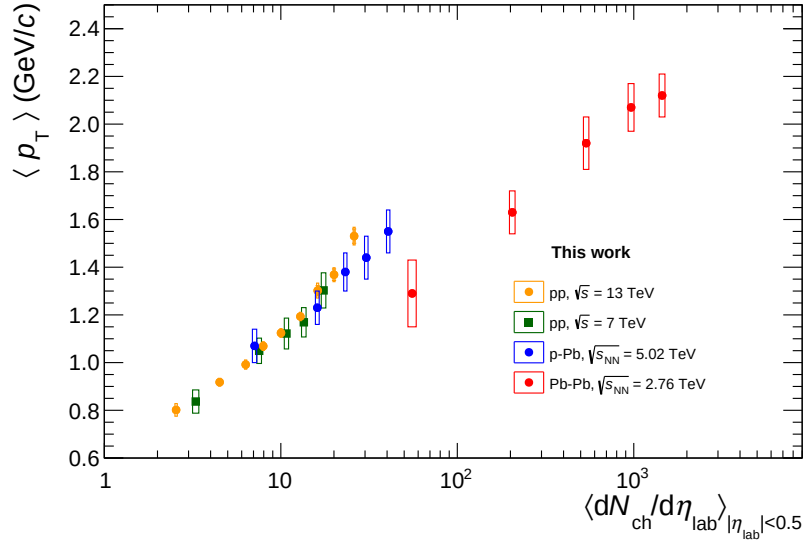
Fig. 5.7.: Transverse-momentum integrated yields dN/dy for deuterons in pp collisions (a) and in all collision systems (b). Data points are measured in pp collisions at $\sqrt{s} = 13$ TeV, pp collisions at $\sqrt{s} = 7$ TeV [85], p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] and in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80]. For pp collisions at $\sqrt{s} = 13$ TeV, the shaded area is the multiplicity-correlated systematic uncertainty.

shown in Fig. 1.11. It is explained as a consequence of the energy loss of partons in a deconfined medium, condition that occurs in Pb–Pb collisions but not in pp and in p–Pb collisions [52], or as a consequence of pressure gradients, which are larger in small systems [53]. From Fig. 5.8b, it comes to light that the kinematics of deuterons works as for the other light-flavoured hadrons.

In the next chapter, the p_T -spectra and the p_T -integrated yields of (anti-)deuterons will be compared with those of protons, in order to constrain the mechanisms of nuclear production, through the comparison with the theoretical predictions that we have at disposal.



(a)



(b)

Fig. 5.8.: Mean transverse momentum $\langle p_T \rangle$ for deuterons in pp collisions (a) and in all collision systems (b). Data points are measured in pp collisions at $\sqrt{s} = 13$ TeV, pp collisions at $\sqrt{s} = 7$ TeV [85], p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] and in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80]. For pp collisions at $\sqrt{s} = 13$ TeV, the shaded area is the multiplicity-correlated systematic uncertainty.

Tab. 5.3.: dN/dy and $\langle p_T \rangle$ for deuterons (a) and anti-deuterons (b), with statistical uncertainties, multiplicity-uncorrelated and multiplicity-correlated systematic uncertainties, reported in the table with this order.

(a) deuterons

Mult. class	$dN/dy (\times 10^{-4})$	$\langle p_T \rangle (\text{GeV}/c)$
I	$15.4 \pm 0.3 \pm 0.4 \pm 1.1$	$1.52 \pm 0.03 \pm 0.03 \pm 0.03$
II	$11.7 \pm 0.1 \pm 0.3 \pm 0.7$	$1.33 \pm 0.03 \pm 0.03 \pm 0.03$
III	$9.1 \pm 0.1 \pm 0.3 \pm 0.6$	$1.30 \pm 0.03 \pm 0.03 \pm 0.03$
IV + V	$7.06 \pm 0.06 \pm 0.22 \pm 0.57$	$1.19 \pm 0.02 \pm 0.01 \pm 0.02$
VI	$5.27 \pm 0.06 \pm 0.25 \pm 0.33$	$1.13 \pm 0.02 \pm 0.02 \pm 0.01$
VII	$4.05 \pm 0.07 \pm 0.25 \pm 0.25$	$1.08 \pm 0.03 \pm 0.02 \pm 0.02$
VIII	$3.05 \pm 0.06 \pm 0.23 \pm 0.19$	$1.02 \pm 0.03 \pm 0.02 \pm 0.02$
IX	$1.85 \pm 0.04 \pm 0.15 \pm 0.11$	$0.91 \pm 0.02 \pm 0.01 \pm 0.02$
X	$0.69 \pm 0.03 \pm 0.09 \pm 0.04$	$0.77 \pm 0.02 \pm 0.04 \pm 0.02$

(b) anti-deuterons

Mult. class	$dN/dy (\times 10^{-4})$	$\langle p_T \rangle (\text{GeV}/c)$
I	$16.0 \pm 0.4 \pm 0.5 \pm 1.8$	$1.57 \pm 0.08 \pm 0.05 \pm 0.03$
II	$12.2 \pm 0.2 \pm 0.4 \pm 1.4$	$1.43 \pm 0.04 \pm 0.04 \pm 0.02$
III	$9.4 \pm 0.1 \pm 0.3 \pm 1.1$	$1.31 \pm 0.03 \pm 0.03 \pm 0.04$
IV + V	$7.13 \pm 0.08 \pm 0.20 \pm 0.79$	$1.21 \pm 0.02 \pm 0.01 \pm 0.03$
VI	$5.34 \pm 0.07 \pm 0.20 \pm 0.59$	$1.12 \pm 0.02 \pm 0.01 \pm 0.03$
VII	$3.99 \pm 0.07 \pm 0.20 \pm 0.44$	$1.06 \pm 0.02 \pm 0.01 \pm 0.03$
VIII	$2.73 \pm 0.04 \pm 0.06 \pm 0.30$	$0.98 \pm 0.01 \pm 0.01 \pm 0.03$
IX	$1.64 \pm 0.03 \pm 0.06 \pm 0.19$	$0.92 \pm 0.01 \pm 0.01 \pm 0.03$
X	$0.59 \pm 0.02 \pm 0.04 \pm 0.07$	$0.82 \pm 0.01 \pm 0.02 \pm 0.02$

B_2 and d/p ratio as a function of multiplicity

In this chapter, the measurements of the coalescence parameter B_2 and of the ratio between the p_T -integrated yields of deuterons and protons will be shown as a function of the event charged particle multiplicity $dN_{\text{ch}}/d\eta$. The observed values will be compared with the predictions of CSM and of the coalescence model. These results together with those presented in previous chapter allow us to have important information to shed light on the nuclei production.

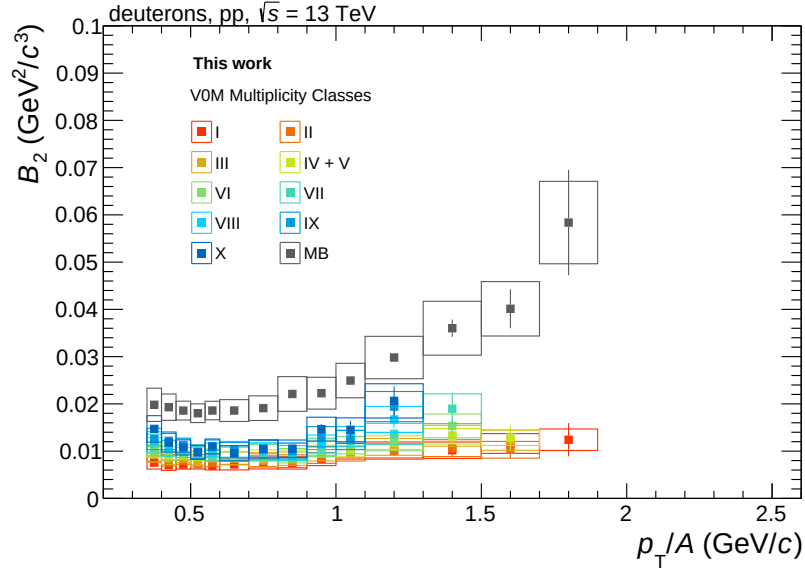
6.1 Coalescence parameter B_2 as a function of p_T

The coalescence parameter B_2 is computed from (anti-)deuteron and proton spectra according to Eq. 2.13. Fig. 6.1 shows B_2 as a function of p_T/A , i.e. the transverse momentum of the proton, for deuterons (Fig. 6.1a) and for anti-deuterons (Fig. 6.1b).

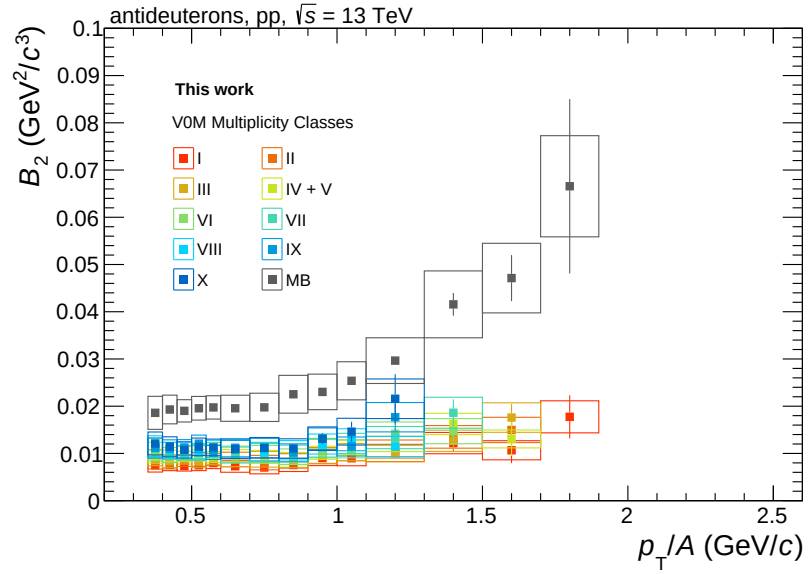
The coalescence parameter is computed separately for deuterons and anti-deuterons, but their values are compatible because the ratio between the production spectra of deuterons and anti-deuterons are compatible with the unity (see Sec. 5.2) and anti-proton spectra are provided as the sum of proton and anti-proton spectra, since their ratio is known to be compatible with the unity.

From the experimental data one can observe that the values of B_2 in different multiplicity classes are very similar to each other, the B_2 computed for the MB analysis shows an offset with respect to the study of the production as a function of the multiplicity. The reason behind this offset is the different normalisation of the spectra. Indeed, as described in Sec. 4.1, MB spectra are normalised to INEL events, while for the analysis as a function of multiplicity spectra are normalised to $\text{INEL} > 0$ events. Since in the B_2 definition the proton spectra appear as squared, the normalisation does not cancel out. For a better visibility, in Fig. 6.2 B_2 is scaled by a power of 2.

In Fig. 6.2 the B_2 reported as a function of the multiplicity shows a slight increase as a function of the transverse momentum. In order to check whether the dependence

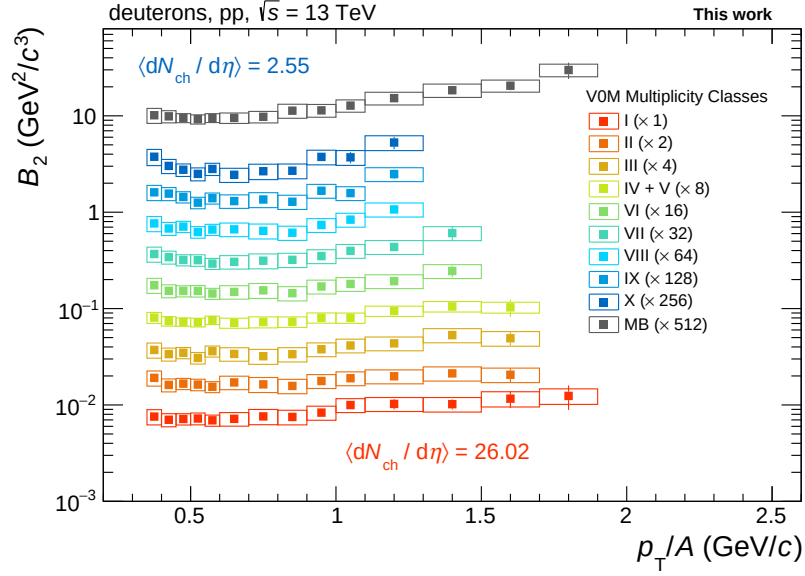


(a) Deuterons

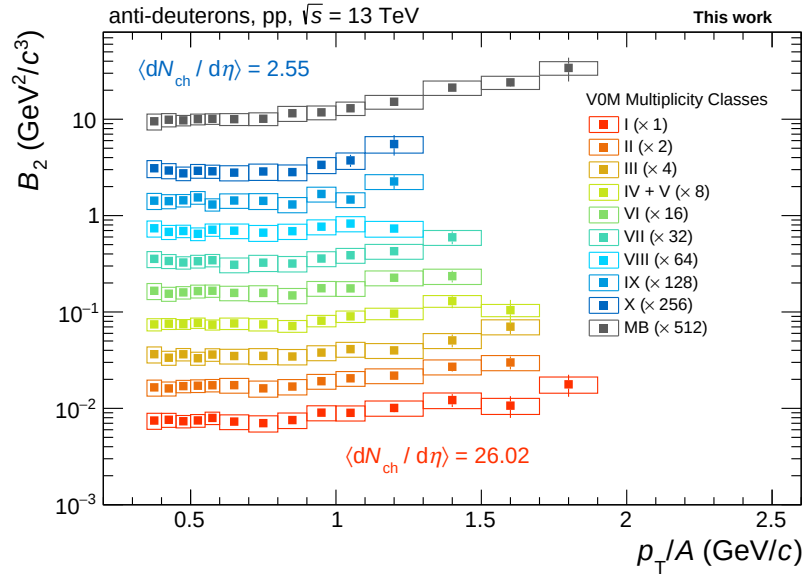


(b) Anti-deuterons

Fig. 6.1.: Coalescence parameter B_2 as a function of p_T/A for all the multiplicity classes under study and for the MB analysis.



(a) Deuterons



(b) Anti-deuterons

Fig. 6.2.: Coalescence parameter B_2 as a function of p_T/A for all the multiplicity classes under study and for the MB analysis, scaled for visibility.

on transverse momentum is significant, B_2 as a function of p_T is fitted with a zeroth order polynomial. The fits are reported in Fig. 6.3, both for deuterons (Fig. 6.3a) and for anti-deuterons (Fig. 6.3b). In every case, the χ^2 of the fit tells us that B_2 as a function of p_T is compatible with a constant. In order to test whether there is an increasing trend with p_T , the uncertainties should be reduced and the measurement should be extended towards higher p_T . In both cases a data sample even larger than the current one would be necessary.

In the MB analysis the increase of B_2 with transverse momentum is much more significant, even if in multiplicity classes it is compatible with a constant. This behaviour has already been observed in p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [85]. In Ref. [85] it is shown that an increasing B_2 can be obtained for the MB analysis or for integrated-multiplicity class even if B_2 is flat in the different multiplicity classes. The cause is the hardening of the proton spectra with increasing multiplicity, that plays a role in the mathematical definition of B_2 .

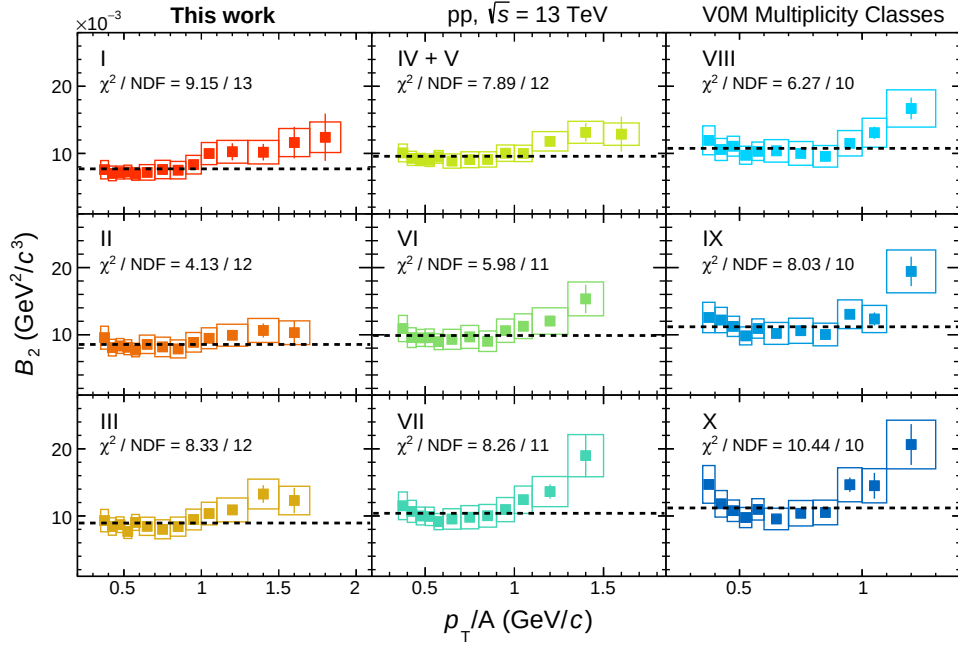
A B_2 flat in p_T is also compatible with the prediction of the most simple coalescence model. B_2 is observed to be compatible with a flat behaviour also in p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] but in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80] it shows a clear increase with p_T (see Fig. 6.4), explained with the presence of radial flow, whose effects are much larger in Pb–Pb collisions than in smaller systems. Moreover, advanced implementations of the coalescence models [34, 38], which take into account space-time distributions of nucleons, could also explain the small raise in p_T , due to the dependence on p_T of the homogeneity radii of the source (see Eq. 2.16).

In the next section, B_2 will be compared across different collision systems and energies and it will be studied as a function of multiplicity.

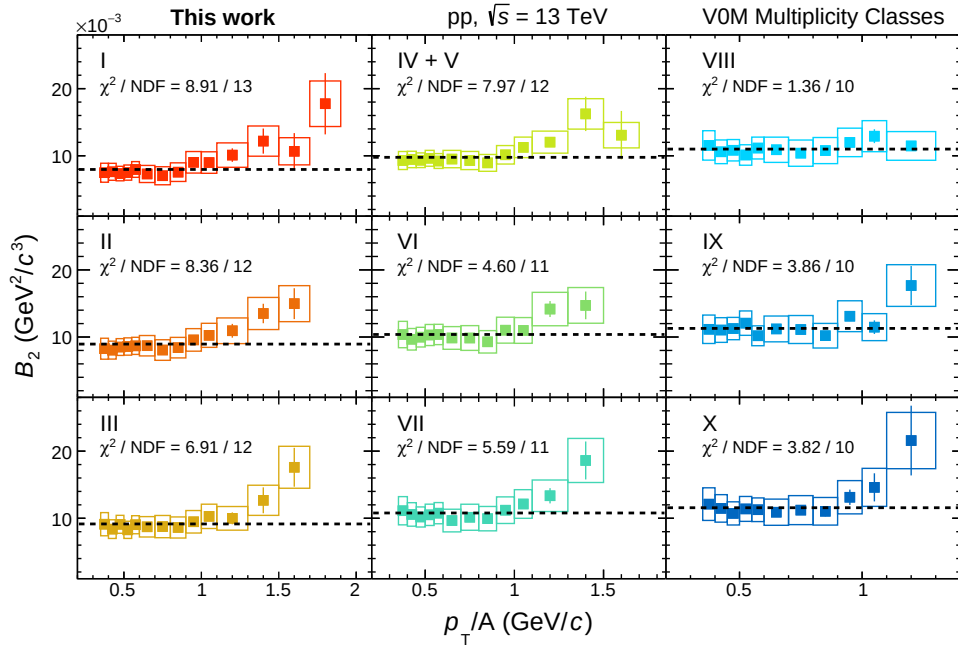
6.2 B_2 as a function of multiplicity

In this section, B_2 is studied as a function of charged particle multiplicity. In order to compare B_2 across different collision systems, B_2 at a fix transverse momentum value, common among all the collision systems, is plotted as a function of charged particle multiplicity. It is not possible to use an average B_2 for each multiplicity class because in Pb–Pb collisions the increase of B_2 with p_T is not negligible.

Fig. 6.5 shows B_2 as a function of multiplicity at $p_T/A = 0.75$ GeV/c. Data corresponds to pp collisions at $\sqrt{s} = 7$ TeV [85] (green points) and at $\sqrt{s} = 13$ TeV (orange points, this work), to p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] (blue

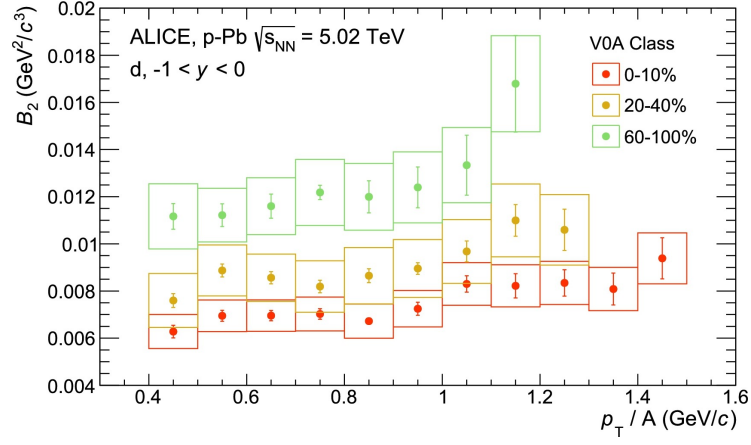


(a) Deuterons

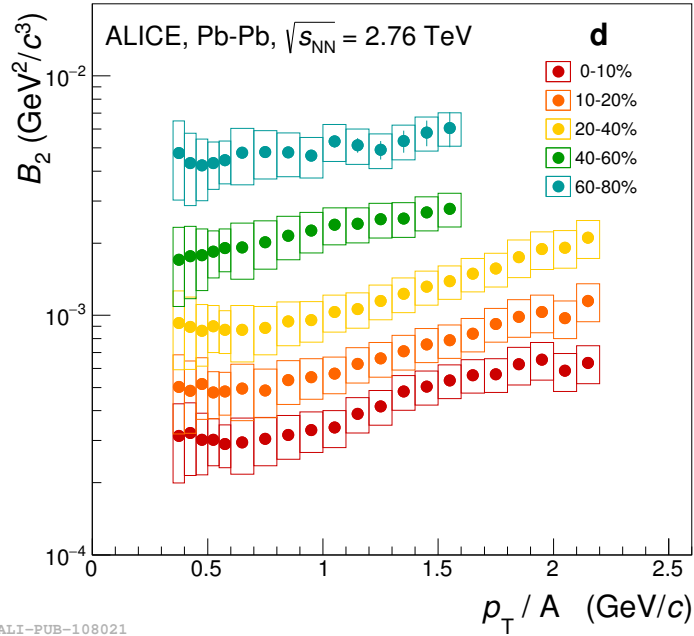


(b) Anti-deuterons

Fig. 6.3.: Fit to the B_2 as a function of p_T with a constant, for deuterons (a) and for anti-deuterons (b), for all the multiplicity classes under study. The fit functions are shown as dashed black lines.



(a) p-Pb, $\sqrt{s_{NN}} = 5.02$ TeV



(b) p-Pb, $\sqrt{s_{NN}} = 2.76$ TeV

Fig. 6.4.: B_2 as a function of p_T in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] (a) and in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80] (b).

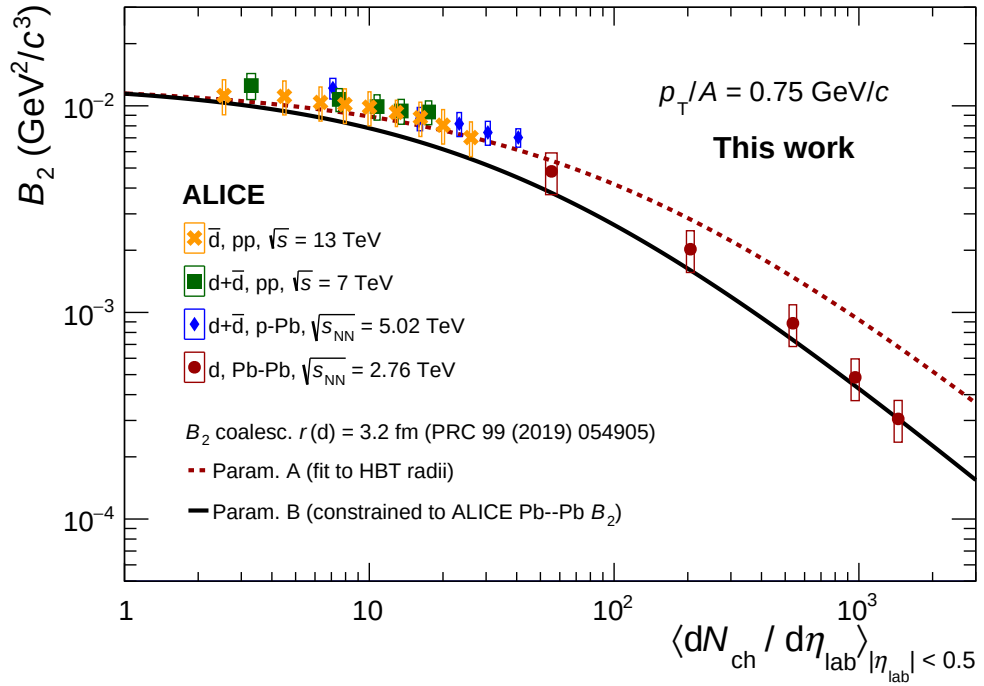
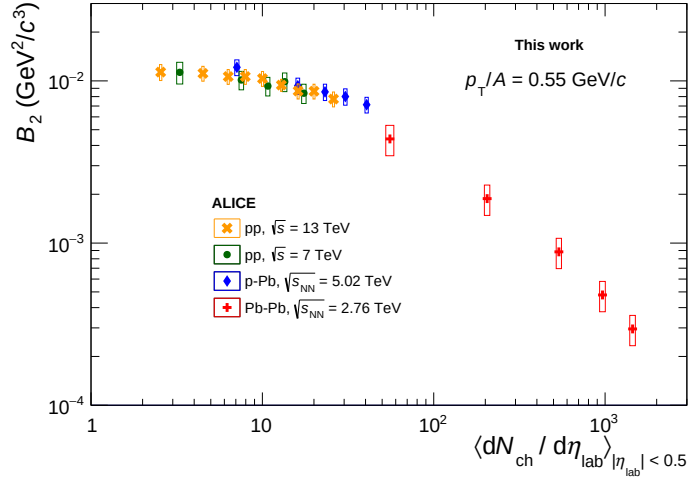
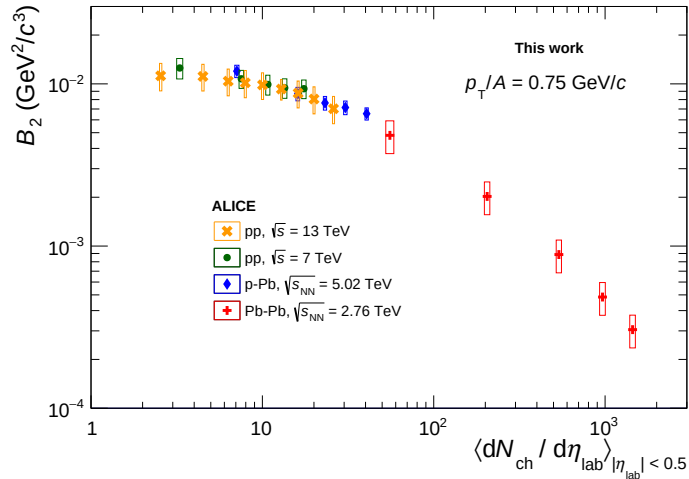


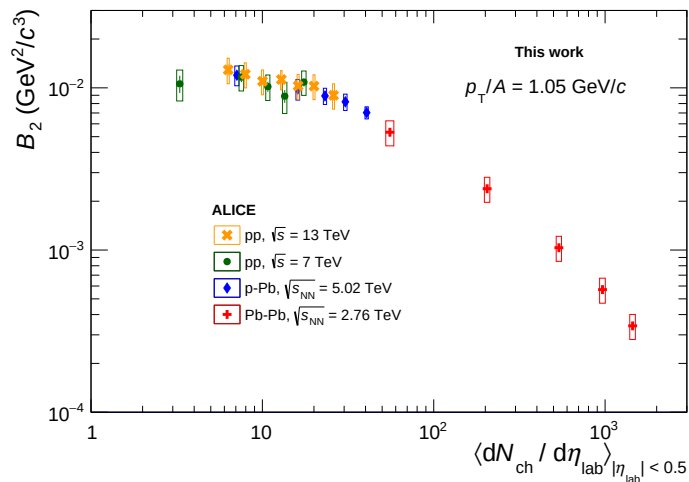
Fig. 6.5.: B_2 evaluated at $p_T/A = 0.75 \text{ GeV}/c$ as a function of $dN_{ch}/d\eta$ for pp collisions at $\sqrt{s} = 7 \text{ TeV}$ [85] and at $\sqrt{s} = 13 \text{ TeV}$ (this work), for p-Pb collisions at $\sqrt{s_{NN}} = 5.02 \text{ TeV}$ [142] and for Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76 \text{ TeV}$ [80]. Data are compared with the predictions of coalescence model for two different parameterisation of the system radius.



(a) $p_T/A = 0.55 \text{ GeV}/c$



(b) $p_T/A = 0.75 \text{ GeV}/c$



(c) $p_T/A = 1.05 \text{ GeV}/c$

Fig. 6.6.: B_2 as a function of charged particle multiplicity at $p_T/A = 0.55 \text{ GeV}/c$ (a), at $p_T/A = 0.75 \text{ GeV}/c$ (b) and at $p_T/A = 1.05 \text{ GeV}/c$ (c).

points) and to Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV [80] (red points). The value of $p_{\text{T}}/A = 0.75$ GeV/ c has been chosen because it is common between all the collision systems, but the behaviour is very similar for the other values of p_{T}/A (see Fig. 6.6).

Thanks to the fine multiplicity classes used in this work it is possible to directly compare data corresponding to pp and p–Pb collisions. The measurements performed in this work cover the whole multiplicity range analysed in pp collisions at $\sqrt{s} = 7$ TeV, extending the measurement both at lower and higher p_{T}/A , and allow for the direct comparison with the measurements obtained in p–Pb collisions, by reaching middle multiplicity classes of the p–Pb measurement.

From the aforementioned comparison it is possible to see that the B_2 measurements in these three different systems and energies are in agreement. Hence, one can conclude that B_2 does not depend neither on the system energy nor on the collision system, but only on the charged particle multiplicity, which reflects the size of the system (see Sec. 1.7).

Considering also higher multiplicities, B_2 shows a smooth evolution passing from small systems to Pb–Pb collisions. This can be interpreted as a hint of a common production mechanism that depends only on the system size and not on the kind or on the energy of the colliding system.

In Fig. 6.5, measurements are compared with the prediction of a coalescence model [34]. The predictions are reported for two different parameterisations of the HBT radius as a function of multiplicity. In the parameterisation corresponding to the solid black line, the relation between HBT radius and multiplicity is fixed in order to reproduce the B_2 measured in the highest multiplicity class in Pb–Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV. In the parameterisation corresponding to the dashed red line, the relation between HBT radius is obtained by fitting the HBT radii measurements as a function of multiplicity (see Sec. 1.7).

B_2 as a function of multiplicity shows two regimes, at low and high multiplicity respectively, that can be explained within the coalescence model. At low multiplicity, B_2 is slightly decreasing. Indeed, in this case the system size is comparable with the deuteron size, expressed as the RMS of the charge wave-function of deuteron (see Sec. 2.2). For this reason, if a(n) (anti-)proton and a(n) (anti-)neutron have a similar momentum, the probability to form a(n) (anti-)deuteron via coalescence is relatively high. According to the definition of B_2 given in Eq. 2.16, the almost-flatness of B_2 is expected, because at low multiplicity, i.e. for a small system radius, the quantum-canonical correction factor $\langle C_d \rangle$ is dominant (see Fig. 2.5).

On the other hand, moving towards higher multiplicity the system size becomes larger and larger than the deuteron size. Therefore, the probability to find a(n) (anti-)proton and a(n) (anti-)neutron with similar momentum in a region with a size compatible with the deuteron one becomes lower and lower. This is reflected by the trend of B_2 in Pb-Pb collisions. As shown in Fig. 2.5, for high multiplicity the quantum correction factor is asymptotically 1 and the value of B_2 is dominated by the system size.

In conclusion, the trend of B_2 as a function of multiplicity can be explained by a coalescence model that takes into account the finite size of the deuteron. However, at the state of arts, neither of the two parameterisations of the HBT radius shown in Fig. 6.5 can describe B_2 both at low and high multiplicity. With new future measurements of the HBT radii as a function of multiplicity, a new and more precise parameterisation could be determined and it could be possible to reproduce data in the whole multiplicity range.

6.3 d/p as a function of multiplicity

The ratio between the p_T -integrated yields of deuterons and protons (d/p) is an important observable to test the theoretical models. Indeed, both coalescence models and SHMs can give predictions about d/p as a function of multiplicity.

In Fig. 6.7, d/p is shown for pp collisions at $\sqrt{s} = 7$ TeV [85] (green points), pp collisions at $\sqrt{s} = 13$ TeV (orange points, this work), p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] (blue points) and Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80] (red points). Measurements are reported as $2 d / (p + \bar{p})$ because proton spectra are provided as the sum of protons and anti-protons and as previously shown deuteron spectra are compatible with those of anti-deuterons.

d/p increases at low multiplicity and saturates at high multiplicity, corresponding to Pb-Pb collisions. As for B_2 as a function of multiplicity (see Sec. 6.2), the transition between pp, p-Pb and Pb-Pb collisions is smooth. This behaviour, as for B_2 , can be interpreted as a hint of a common production mechanism that depends only on the system size and not on the kind and energy of the system.

Fig. 6.7b focuses on the low multiplicity ($0 < \langle dN_{ch}/d\eta \rangle < 30$) reporting pp collisions at $\sqrt{s} = 7$ TeV and at $\sqrt{s} = 13$ TeV and p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. As for B_2 , the measurement in pp collisions at $\sqrt{s} = 13$ TeV covers the full multiplicity range of pp collisions at $\sqrt{s} = 7$ TeV, extending it at low and high multiplicity, and reaches the intermediate multiplicity classes analysed in p-Pb collisions. The

three sets of measurements show an excellent agreement, further confirming that the production mechanisms of nuclei depend only on the system size.

The measurements reported in Fig. 6.7 are also compared with the predictions of the coalescence model [38] (solid magenta line) and of Thermal-FIST CSM [73] (shaded grey area). For CSM, the shaded grey area corresponds to the predictions obtained for a chemical freeze-out temperature $T_{ch} = 155$ MeV and a correlation volume V_C spanning from 1 (solid black line) to 3 (dashed black line) rapidity unities (see Sec. 2.1.2).

At low multiplicity d/p increases, as described by both the coalescence model and the CSM. For CSM, data suggest a correlation volume V_C that is intermediate between the limits shown in Fig. 6.7, i.e. $dV/dy < V_C < 3 dV/dy$. According to the coalescence model, in this multiplicity region in which the system size is comparable with the deuteron size, the raise with multiplicity is explained in terms of the interplay between the system size and the proton yield. Indeed, both the system size and the proton yield increase with multiplicity, but at a different rate. Therefore, as long as the system size is smaller than the deuteron size, if a(n) (anti-)proton and a(n) (anti-)neutron are produced and they have similar momenta, they are close enough to form a(n) (anti-)deuteron via coalescence. Hence, in this multiplicity region d/p is dominated by the proton yield. On the other hand, according to the CSM the raise with multiplicity is explained in terms of canonical suppression (see Sec. 2.1.2).

At high multiplicity, corresponding to Pb–Pb collisions, d/p saturates. The Pb–Pb measurement in the highest multiplicity classes seem to suggest a possible decrease, but with the current uncertainties the hypothesis of flatness cannot be rejected. Both the coalescence model and CSM predict this saturation at high p_T . In the frame of coalescence model, this is due to the fact that the system size is larger than the deuteron size. Thus, even if the nucleon yield increases, it becomes less likely to find a proton and a neutron close to each other to form a nucleus via coalescence. Instead, in the frame of SHM, with increasing multiplicity the canonical description of the system tends to the grand-canonical one, where d/p is expected not to depend on multiplicity. Both the models fail in describing the hypothetical decrease at very high multiplicity.

Both the coalescence model and CSM can describe the trend of d/p as a function of multiplicity. However, canonical suppression in CSM cannot reproduce the ratio between the integrated yields of protons and pions (p/π), which have the same baryon-content ratio as d/p [146]. Fig. 6.8, taken from Ref. [146], shows the ratios p/π , K/π , ϕ/π , Λ/π , Ξ/π and Ω/π . The data are measured by ALICE in pp collisions at $\sqrt{s} = 7$ TeV, p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV and Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV and are compared with the CSM predictions, for a chemical

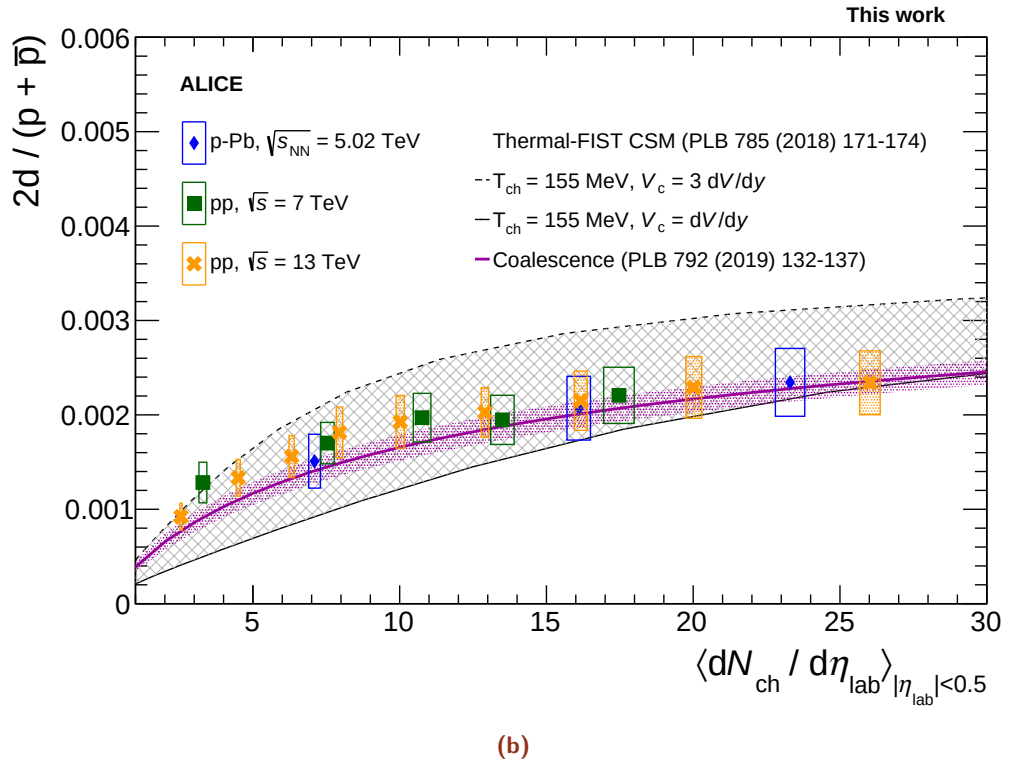
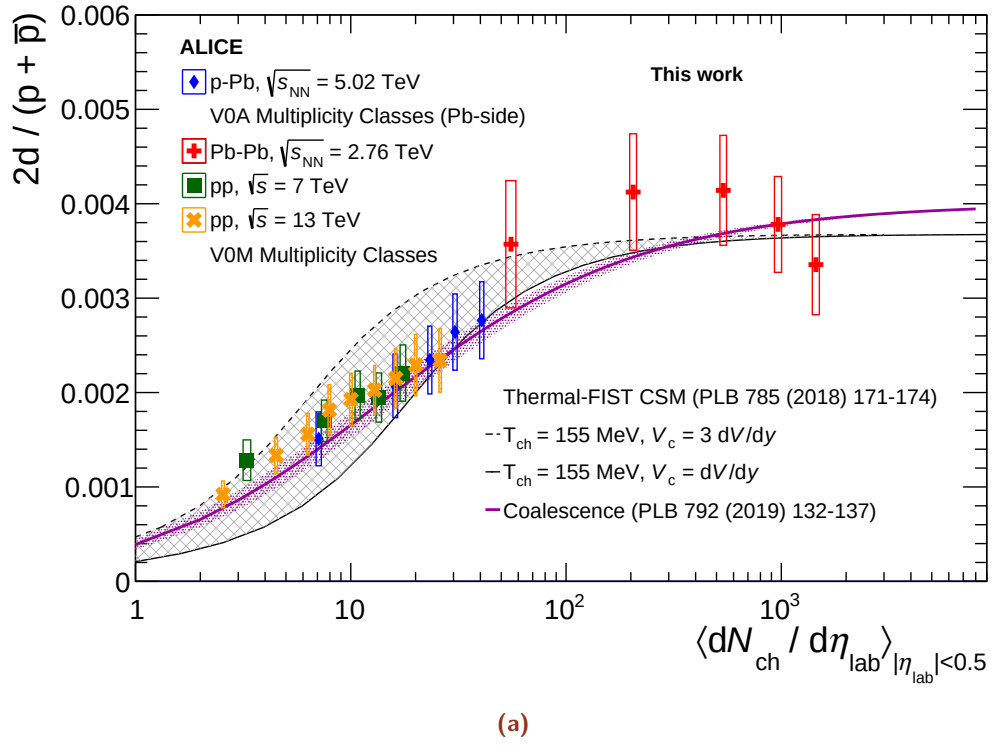


Fig. 6.7.: Ratio between the p_T -integrated yields of deuterons and protons d/p as a function of $dN_{ch}/d\eta$ for pp collisions at $\sqrt{s} = 7$ TeV [85] and at $\sqrt{s} = 13$ TeV (this work), for p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] and for Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80]. For pp collisions at $\sqrt{s} = 13$ TeV, the shaded area represents the multiplicity-correlated systematic uncertainty. Data are compared with the predictions of coalescence model and with the prediction of CSM with two different values of correlation volume. Panel (b) is the same measurement shown in panel (a), but focused on low multiplicity.

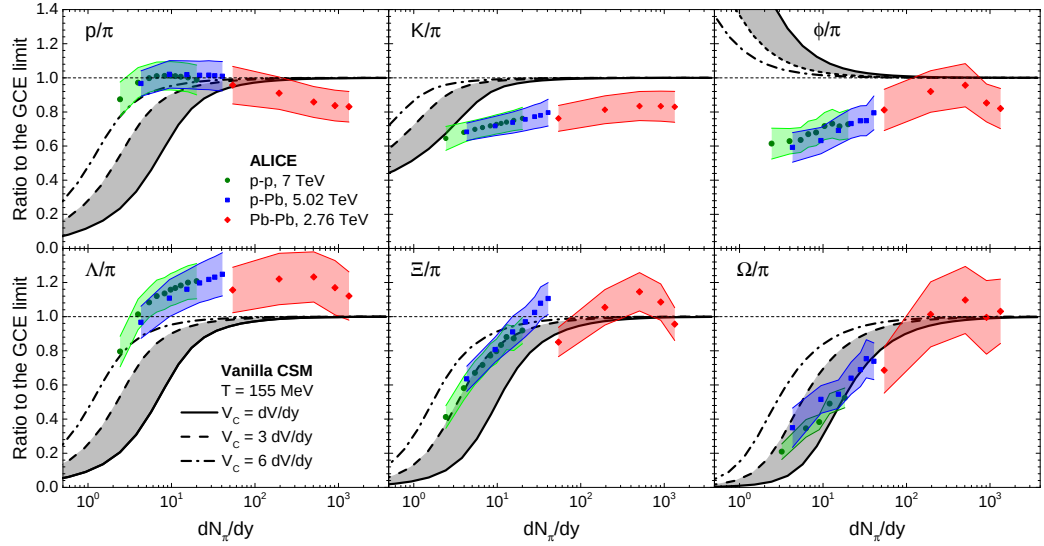


Fig. 6.8.: Ratio between p_T -integrated yields of light-flavoured hadrons and pions compared with the prediction of CSM. Both calculated results and the data are scaled by the grand-canonical limiting values as evaluated in the CSM with $T_{ch} = 155\text{MeV}$ and for $\mu_B = 0$. Figure taken from Ref. [146].

freeze-out temperature $T = 155\text{ MeV}$ and for a correlation volume spanning from 1 to 6 rapidity units. Both the data and the model predictions are scaled by the gran canonical (GCE) limiting valued evaluated at $T = 155\text{ MeV}$ and $\mu_B = 0$. The CSM describes well Ξ/π and Ω/π and reproduces the trend observed for Λ/π , even if it slightly undershoots the data. The CSM also overshoots K/π and it does not reproduce the served trend at low multiplicity. The predictions for ϕ/π is opposite to the behaviour in data. The description of the ratios to pions for the hadron species containing strange quarks is improved by adding a strangeness saturation factor γ_S [146]. Finally, the CSM does not describe the measured p/π with a correlation volume $dV/dy < V_c < 3dV/dy$ as used for d/p . Indeed, the data do not show a conclusive evidence for a suppression of p/π at low multiplicity. Although the net baryon content and the net strangeness content are the same, the p/π ratio requires a very large correlation volume (6 rapidity units), while the d/p ratio requires a correlation volume of 1-3 rapidity units. This measurement disfavours a statistical-thermal production of nuclei in small systems.

Summary and conclusions

The goal of this thesis project was the study of (anti-)deuteron production in pp collisions at $\sqrt{s} = 13$ TeV. This study has been carried out on the data collected by ALICE in the years 2016 and 2017, consisting of about 1 billion minimum bias (MB) pp collisions provided by the LHC. Thanks to the large data sample, the largest ever used in a nuclei-related analysis within the ALICE collaboration, it has been possible to study the production of (anti-)deuterons in nine fine multiplicity classes and to extend the p_T reach of the analysis with respect to the previous measurement performed with the ALICE data in pp collisions at $\sqrt{s} = 7$ TeV. In particular, it has been possible to extend the multiplicity coverage of the analysis in pp collisions at $\sqrt{s} = 7$ TeV, both at low and at high multiplicities, thus reaching the multiplicity range of p–Pb collisions. For this reason, this work represents a step forward in the comparison of the production of light (anti-)nuclei across different collisions systems, such as pp and p–Pb collisions, at similar multiplicities and then in the understanding of the production mechanism of light (anti-)nuclei at the LHC energies.

Following the procedure described in Chapter 4, the transverse momentum spectra of deuterons and anti-deuterons have been measured (Fig. 5.1). The ratio between the two (Fig. 5.3) is compatible with the unity and therefore the same production of deuterons and anti-deuterons is observed, as expected at the LHC energies, where the baryon-chemical potential is null. From the p_T -spectra, the p_T -integrated yield dN/dy and the mean transverse momentum $\langle p_T \rangle$ are evaluated. This is the first nuclei-related analysis by the ALICE Collaboration in which the multiplicity correlated systematic uncertainty on those quantities is evaluated and the systematic uncertainties are shown to be in large part correlated. Both dN/dy (Fig. 5.7) and $\langle p_T \rangle$ (Fig. 5.8) are in agreement with the measurements in pp collisions at $\sqrt{s} = 7$ TeV and in p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV at similar multiplicities. On the one hand, dN/dy smoothly as a function of multiplicity from pp to Pb–Pb collisions, as observed for the other light-flavoured hadrons, confirming that hadron chemistry is driven only by the event multiplicity. On the other hand, $\langle p_T \rangle$ evolves smoothly with multiplicity in pp and p–Pb collisions, but there is a discontinuity passing to Pb–Pb collision. Also this aspect is in common with the $\langle p_T \rangle$ of the other light-flavoured hadrons and it is explained as a consequence of the energy loss of partons in a deconfined medium, which is produced in Pb–Pb collisions but not in smaller systems.

The main results of this work are the measurements of the coalescence parameter B_2 (Eq. 2.13) and of the ratio between the p_T -integrated yields of deuterons and protons d/p , because they are instrumental in the direct comparison with the predictions of the theoretical models. These two observables are shown in Fig 7.1 and in Fig 7.2, respectively, as a function of multiplicity. Both these observables are obtained from the comparison of the production of deuterons with that of protons. These measurements are compared with the predictions of two classes of theoretical models, namely the coalescence model (Sec. 2.2) and the statistical hadronisation model (SHM) (Sec. 2.1). This is the first deuteron-related analysis in which these observable are compared with the predictions of theoretical models. According to the coalescence model, nucleons that are close in phase-space can form a nucleus via coalescence. B_2 is the key observable of the model, because it is related to the probability to form a nucleus via coalescence. Simple implementations of the coalescence model do not take into account the space distribution of nucleons, but only their momentum distribution. According to the simple coalescence model, B_2 does not depend on p_T or on multiplicity. More advanced implementations of the coalescence model take into account the system size, which is related to the charged particle multiplicity (Sec. 1.7). In particular, these implementations are based on the superposition of the (anti-)deuteron wave-function with the phase space of the nucleons. Due to the interplay between system and deuteron size, B_2 is expected to depend on the system size, hence on the event multiplicity (Eq. 2.16). The SHM, instead, can predict the abundances of hadrons, including (anti-)nuclei, based on the assumption of statistical equilibrium at the chemical freeze-out, when the abundances of stable particles are fixed. In Pb–Pb collisions, characterised, by a large system size, a grand-canonical approach is used (Sec. 2.1.1). In a grand-canonical description ratios between particle yields, such as d/p , are expected not to depend on the event multiplicity. On the other hand, in pp and p–Pb collisions, the condition for a grand canonical description are not satisfied and a canonical description is used (Sec. 2.1.2). In the canonical statistical hadronisation model (CSM), the exact conservation of baryon number B , strangeness content S and electric charge Q is imposed. In particular, $B = S = Q = 0$. According to CSM, d/p is expected to increase with the system size, hence with multiplicity, for small systems. For high multiplicities, the predictions of the canonical approach tend to those of the grand-canonical description.

In this work B_2 as a function of p_T has been measured in multiplicity classes and for the minimum bias (MB) analysis. In the analysis in multiplicity classes, B_2 is compatible with the hypothesis of a flat behaviour (Fig. 6.3), in agreement even with the predictions of a simple coalescence model, which does not take into account space distribution of nucleons. However, a hint of increase at high p_T could be present, even if the hypothesis of flatness cannot be rejected. This small raise could be explained in the frame of an advanced coalescence model as the consequence

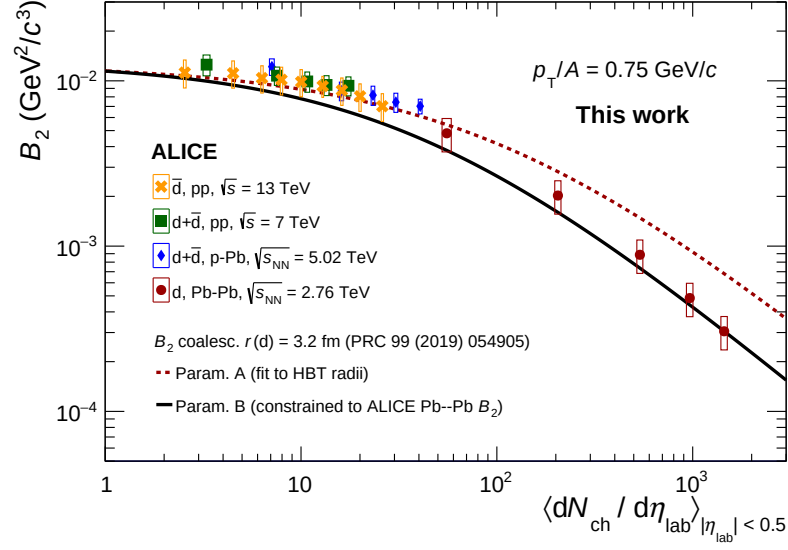


Fig. 7.1.: B_2 evaluated at $p_T/A = 0.75$ GeV/c as a function of $dN_{ch}/d\eta$ for pp collisions at $\sqrt{s} = 7$ TeV [85] and at $\sqrt{s} = 13$ TeV (this work), for p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV [142] and for Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [80]. Data are compared with the predictions of coalescence model for two different parameterisation of the system radius (same figure as Fig.6.5).

of the dependence of the homogeneity volume on p_T . The B_2 for the MB analysis, instead, increases with p_T . This is expected and it is caused by the hardening of proton spectra and its effect in the mathematical definition of B_2 (Eq. 2.13). This effect was already observed in previous analyses.

The coalescence parameter B_2 at fixed p_T/A shows a smooth evolution as a function of multiplicity (Fig. 6.5 or Fig. 7.1). The measurements in pp at $\sqrt{s} = 13$ TeV (this work) are in agreement with those of pp collisions at $\sqrt{s} = 7$ TeV and with p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV at similar multiplicities. Two regimes of B_2 can be distinguished, corresponding to two multiplicity regions. At low multiplicity, corresponding to small systems, B_2 is slightly decreasing. This behaviour is in agreement with the predictions of the coalescence model. Indeed, at low multiplicity the system size, expressed as the Hanbury-Brown-Twiss radius, is comparable with the deuteron size and therefore it is more likely to produce (anti-)deuterons via coalescence. Moving towards higher multiplicities, the system size becomes larger and larger than the deuteron size and the probability to form a(n) (anti-)deuteron via coalescence decreases: even if more nucleons are produced, they could be so far not to be able to form a nucleus via coalescence. The data are compared with the predictions of the coalescence model for two different parameterisations of the HBT radius as a function of multiplicity. In the first parameterisation, the relation between HBT radius and multiplicity is fixed in order to reproduce the B_2 measured in the highest multiplicity class in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. In the second

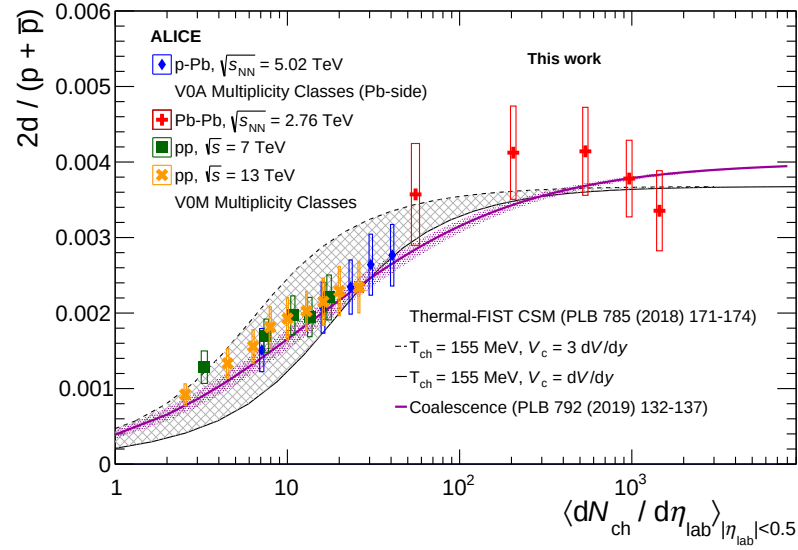


Fig. 7.2.: d/p as a function of $dN_{ch}/d\eta$ (same figure as Fig.6.7).

parameterisation, the same relation is obtained by fitting the HBT measurements as a function of multiplicity. Even if the two predictions describe quite well the observed trend, neither of them reproduces the data over the whole multiplicity range.

The coalescence model describes quite well the trend of B_2 as a function of multiplicity, but it does not reproduce the data over all the multiplicity range with either of the two parameterisations of the HBT radius as a function of multiplicity, one obtained by fitting the HBT measurements as a function of multiplicity, the other by fixing the parameter of the relation between HBT radius and multiplicity to reproduce B_2 in the most central Pb–Pb collisions.

In the parameterisation corresponding to the solid black line, the relation between HBT radius and multiplicity is fixed in order to reproduce the B_2 measured in the highest multiplicity class in Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. On the other hand, in the parameterisation corresponding to the dashed red line, the relation between HBT radius is obtained by fitting the HBT radii measurements as a function of multiplicity (see Sec. 1.7).

As for B_2 , also d/p shows a smooth evolution as a function of multiplicity (Fig. 6.7 or Fig. 7.2) and the measurements in pp at $\sqrt{s} = 13$ TeV are in agreement with those of pp collisions at $\sqrt{s} = 7$ TeV and with p–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV at similar multiplicities. Also in this case two regimes, corresponding to the high and the low multiplicity regions respectively, can be distinguished. At low multiplicities, where the system size is comparable with the deuteron size, d/p increases. In the frame of

coalescence model, the increase of d/p is explained with the different increase of the system size and of the proton yield as a function of multiplicity. Indeed, as long as the system size is smaller than the deuteron size, if a(n) (anti-)proton and a(n) (anti-)neutron are produced and they have similar momenta, they are close enough to form a(n) (anti-)deuteron via coalescence. Hence, in this multiplicity region d/p is dominated by the proton yield. According to CSM, instead, the raise is a consequence of the exact conservation of the baryon number B (canonical suppression). At high multiplicity, d/p saturates. In the frame of coalescence model, this is due to the fact that the system size is larger than the deuteron size. Therefore, even if the nucleon yield increases, it becomes less likely to find a proton and a neutron with similar momenta and close enough to form a nucleus via coalescence. According to the SHM, with increasing multiplicity the canonical description of the system tends to the grand-canonical one, where d/p is expected not to depend on multiplicity. At very high multiplicity, data points seem to suggest a hint of decrease, but at present the uncertainties are too large to reject the hypothesis of a flat behaviour. If it were really present, both models would fail to reproduce the experimental trend of the data. The coalescence model can describe d/p over the whole multiplicity range. Also CSM can reproduce the data, providing a correlation volume $dV/dy < V_C < 3 dV/dy$. However, CSM fails in the description of the ratio between the p_T -integrated yields of protons and pions p/π : it cannot describe with the same correlation volume and temperature d/p and the p/π simultaneously, even though they are characterised by the same net baryon content and the same net strangeness content.

In conclusion, B_2 and d/p show a smooth evolution with multiplicity, regardless of the collision system and energy. This behaviour suggests the presence of a common production mechanism for (anti-)deuterons that depends only on the system size. Both the coalescence model and the CSM qualitatively describe the trend observed for B_2 and of d/p as a function of multiplicity. However, CSM cannot describe d/p and p/π simultaneously with the same correlation volume and temperature and can be excluded. At the state of arts, the experimental measurements have reached a good precision and a good multiplicity coverage and some improvements in the theoretical models are needed for a better comparison.

Appendices

Cluster encoding for the new ALICE ITS

In this appendix, the development of the algorithm for the encoding of the clusters of pixels of the new ALICE ITS will be described. The first version of the software was implemented by me and was the subject of my Master thesis. The further implementation and the porting of the code to O², the framework for online and offline processing that will be used in ALICE during LHC Run 3, represented my Service Task for the ALICE Collaboration.

A.1 Upgrade of the Inner Tracking System

In LHC Run 3 period of data acquisition (2021), the luminosity of the collider will be higher than in LHC Run 2. For Pb–Pb collisions, the interaction rate will increase from the current value of 8 kHz to 50 kHz, corresponding to an instantaneous luminosity of $6 \mu\text{b}^{-1}\text{s}^{-1}$. For Run 3, the ALICE physics program includes the measurement of rare processes, like the production of Λ_c and Λ_b baryons and the production of hyper-nuclei heavier than ${}^3_\Lambda\text{H}$, which is currently the only hyper-nucleus observed in heavy-ion collisions. For these goals, both a high integrated luminosity and a high resolution on the reconstruction of secondary vertices is needed.

In order to achieve the physics goals for Run 3, the ALICE apparatus is currently undergoing an upgrade. In particular, the current Inner Tracking System (ITS, see Sec. 3.2.1) will be completely replaced by a new ITS, which is now under commissioning. Information about the new ITS can be found in Ref. [147].

The new ITS will consist of seven layers of silicon pixel detectors. The layout of the new ITS is shown in Fig. A.1a and its geometrical parameters are reported in Tab. A.1. The first additional layer is closer to the collision point than the current one, thanks to the installation of a new 0.8 mm thick beryllium beam pipe, with a diameter of 19.2 mm. The closer first layer improves the resolution on the impact parameters by a factor of 3 and of 5 in transverse ($r\varphi$) and longitudinal (z) directions at $p_T = 1 \text{ GeV}/c$, respectively. As a consequence of the proximity to the collision point, the expected radiation load for the first layer is expected to be 700 kRad of

Tab. A.1.: Geometrical parameters of the new ITS.

Parameters	L0	L1	L2	L3	L4	L5	L6
Min. radial position (μm)	22.4	30.1	37.8	194.4	243.9	342.3	391.8
Max. radial position (μm)	26.7	34.6	42.1	197.7	247.0	345.4	394.9
Length (mm)	271	271	271	843	843	1475	1475
Pseudorapidity coverage	± 2.5	± 2.3	± 2.0	± 1.5	± 1.4	± 1.4	± 1.3
Nr. pixel chips	108	144	180	2688	3360	8232	9408

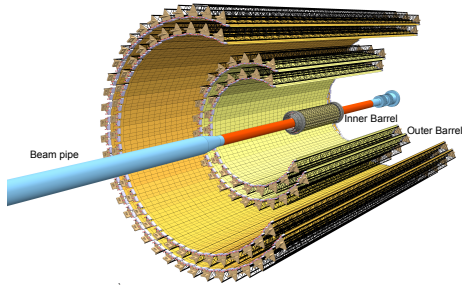
Tab. A.2.: Pixel-chip requirements for the construction of the new ALICE ITS from the Technical Design Report (TDR) [147] (second and third columns) and current performance of ALPIDE sensors (last column).

Parameters	Inner Barrel	Outer Barrel	ALPIDE
Silicon thickness (μm)	50	100	✓
Spatial resolution (μm)	5	10	$\simeq 5$
Chip dimension (mm^2)	15×30		✓
Power density	$< 300 \text{ mW/cm}^2$	$< 100 \text{ mW/cm}^2$	✓
Integration time (μs)	< 30		< 2
Efficiency	$> 99 \%$		
Fake-hit rate ($\text{event}^{-1} \text{ pixel}^{-1}$)	$< 10^{-6}$		$\ll 10^{-6}$
NIEL radiation tolerance ($1 \text{ MeV n}_{\text{eq}}/\text{cm}^2$)	1.7×10^{13}	1.0×10^{13}	✓
TID radiation tolerance (krad)	2700	100	tested at 350

Total Ionising Dose (TID) and $1 \text{ MeV n}_{\text{eq}}/\text{cm}^2$ of Non Ionising Energy Loss (NIEL), including a factor of 10.

The geometry and the finer segmentation of the new ITS will improve both the tracking efficiency and the resolution on the impact parameters. The seven layers can be divided into Inner Barrel, consisting of the three innermost layers, and Outer Barrel, consisting of the four outermost layers. All the layers are equipped with Monolithic Active Pixel Sensors (MAPS), with a pixel size of $O(30 \times 30 \mu\text{m}^2)$, to be compared with the size of the current pixel size of the SPD, corresponding to $50 \times 425 \mu\text{m}^2$. The total sensitive area is about 10 m^2 , with 12 billion pixels with binary readout. The characteristics of the pixel sensors, called ALice Pixel Detector (ALPIDE), can be found in Tab. A.2.

Thanks to the reduced material budget, the multiple scattering for low-momentum particles is reduced, improving tracking efficiency and impact parameter resolution. The low material budget is obtained with a $50 \mu\text{m}$ ($100 \mu\text{m}$) thick sensors in the



(a) Layout of the new ITS.



(b) Three half barrels of the outermost layers.

Fig. A.1.: Layout of the new ITS (a) and three half barrels of the outermost layers (b).

Inner (Outer) Barrel. Moreover, optimised power consumption of the front-end electronics allows for room temperature operations, with a lightweight cooling system. Finally, the support structure is made of ultra light carbon fibre. The total material budget is 0.3% X_0 for the Inner Barrel layers and 1.0% X_0 for the Outer Barrel ones.

The readout rate will be greatly improved. The current ITS has a 1 kHz readout limits, not sufficient collect all the Pb–Pb collisions delivered by the LHC (8 kHz). With the new ITS it will be possible to read Pb–Pb collisions at the peak rate of 50 kHz and pp collisions at 400 kHz, with the possibility to read data continuously.

A.2 Upgrade of the Online-Offline systems

During Run 3 the data throughput of the ALICE experiment will reach 1.1 TB/s [148], mostly coming from the TPC (1 TB/s). The new ITS will have a throughput of 40 GB/s, larger than the value for the whole current ALICE apparatus, i.e. 25 GB/s. The data rate is too large to be written unprocessed and uncompressed. Therefore, most of Offline processes, in particular reconstruction and data calibration, will be moved online in a dedicated facility. This project is called O^2 and its goal is to reduce the time between the data acquisition and the data analysis, providing event building, reconstruction, calibration, online and quasi-online selection.

The data handling paradigm will be changed by the continuous readout for the detectors that supports it. Instead of triggered events, data will be read out and processed in Time Frames (TF), which are groups of readout snapshot of the detector. The O^2 infrastructure will be a High Performance Computing (HPC) facility hosting heterogeneous hardware, including 250 First Level Processor worker nodes equipped

with Field Programmable Gate Arrays (FPGA) and 1500 Event Processing Nodes equipped with Graphics Processing Units (GPU).

As shown in Ref. [148], the data need to be compressed by a factor of 20 before their storage. The idea behind this work is that the ITS data format, i.e. the cluster of pixels, contains redundant information. By optimising the format and by compressing the data with the Huffman coding, the ITS data rate will be reduced to ~ 26 GB/s or to ~ 4.3 GB/s, depending on the noise occupancy. More details about the data format and compression will be given in Sec. A.3.

A.3 Cluster data format

The new ITS is characterised by a binary readout: it can provide only space information about the crossing point of a charged particle within the active area of the sensor. When a charged particle crosses the sensor, it creates some electron-hole pairs. Electrons are collected by a dedicated diode and if the number of collected electrons exceeds a threshold the pixel delivers a signal. Signal pixels are called *firing pixels*. A cluster of pixels (hereinafter called simply cluster) is a set of adjacent firing pixels. For the track reconstruction, the only needed piece of information is an impact point position, with its uncertainty.

Each cluster can be completely described with two geometrical attributes:

1. the *topology*, which is the pattern of firing pixels;
2. the position of the cluster within the sensor.

The topology includes few pieces of information, which are common to all the cluster with that particular topology:

1. the *bounding box*, which is the smallest rectangle that can entirely contain the cluster;
2. the number of firing pixels;
3. the *Centre Of Gravity* (COG) position within the bounding box.

The position of the cluster within the sensor is expressed in terms of the position of a reference pixel within the sensor matrix, as a (row, column) coordinate. The reference pixel is the top-left corner pixel within the bounding box. The position of

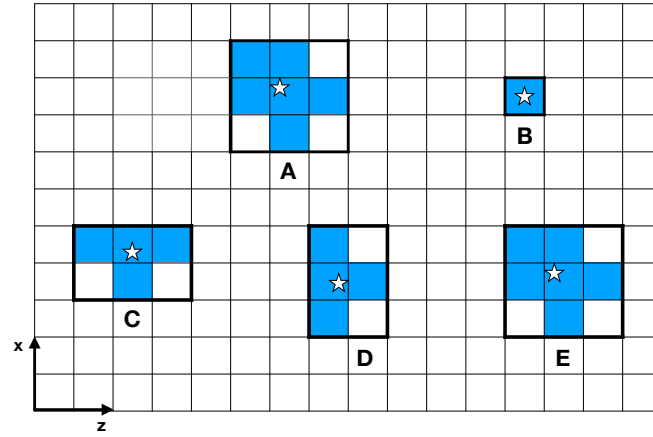


Fig. A.2.: Some examples of cluster of pixels. The thick black line and the white star are the bounding box and the Centre Of Gravity (COG) of the cluster. Clusters A and E have the same topology. Clusters C and D have different topologies, even if one can be transformed into the other through a rotation. B is a single-pixel cluster.

the COG within the bounding box is also expressed relatively to the centre of the reference pixel.

Generally, in clusters with the same topology the true impact point of the particle is different. The COG is the best estimation of the impact point. Its uncertainty is taken from a uniform distribution within a range equal to the size of the cluster. Being x and z the directions of the sensor reference frame (see Fig. A.2), the uncertainties on the impact point position are:

$$\sigma_x = \frac{N_{rows} \cdot d_x}{\sqrt{12}} \quad \sigma_z = \frac{N_{columns} \cdot d_z}{\sqrt{12}} \quad (\text{A.1})$$

where d_x and d_y are the pitch of the pixel along x and z directions, respectively. N_{rows} and $N_{columns}$ are the number of rows and columns of the bounding box, respectively. Alternatively, the uncertainty on the impact point can be obtained from a MC simulation. Indeed, for all the clusters with the same topology, a distribution of the distance between the true impact point and the COG can be obtained. The RMS of this distribution can be used as uncertainty.

The cluster topology and the cluster position within the sensor are all the information needed to evaluate the impact point and its uncertainties. For this reason, a cluster can be completely described with two attributes: an identifier of its topology and the coordinates of its reference pixel. The aim of this work is the creation of a dictionary

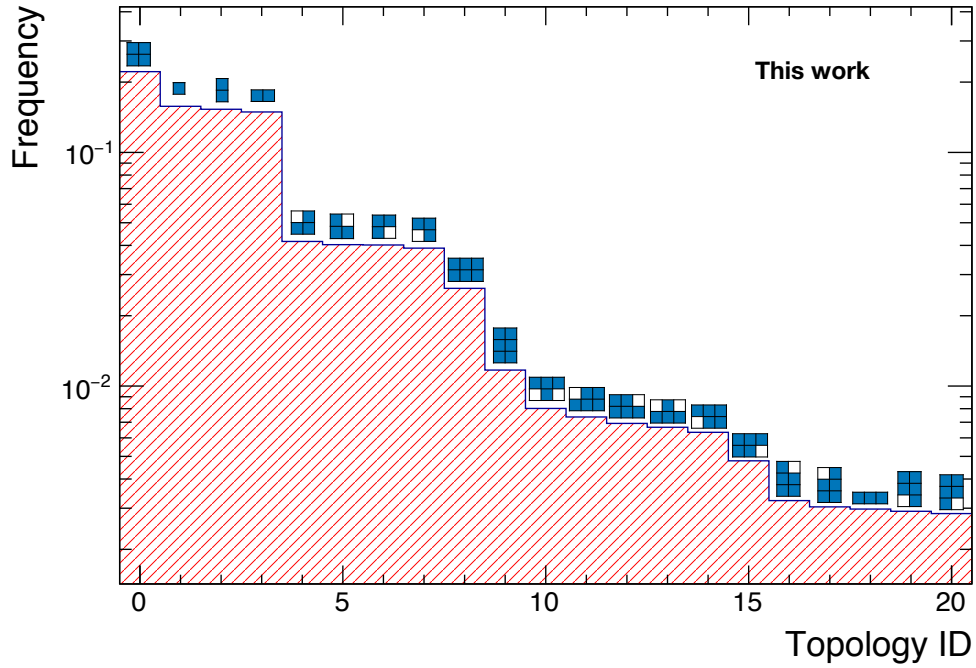


Fig. A.3.: Cluster topology distributions. The twenty most common topologies are shown and they are sorted in descending order of frequency. The distribution has been obtained simulating 50 pp collisions.

containing an entry for each topology. In this way, the topology can be stored as a reference to the corresponding entry in the dictionary.

Fig. A.3 shows the topology distribution obtained with 50 simulated pp collisions. This distribution is highly not uniform and this characteristic can be exploited to compress topology information using Huffman coding [149]. The Huffman coding is a lossless data compression algorithm, which provides as output a variable-length code table for encoding a source symbol, like, for example, the character in a file. The table is derived from the frequency of occurrence of each possible value of a source symbols. Using Huffman coding, more common symbols are generally represented using fewer bits than less common symbols. In this case of use, the symbol source is the topology identifier (ID). The highly non uniformity of the distribution guarantees a good compression of the source, since the most common topology ID is mapped into the shortest bit string.

The compression is obtained at the cost of creating a Huffman table and consulting it during the coding and the decoding processes, operations which are time consuming. For this reason, since the compression/decompression time is an important factor, it is necessary to pay attention to the size the table, because the look-up time is

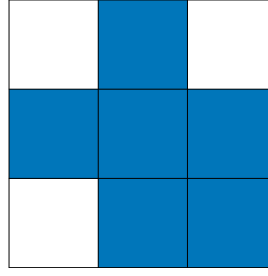


Fig. A.4.: Example of cluster topology. The bounding box is 3×3 . The raw format (bit sequence) of the topology is 00000011 00000011 01011101 1000000. The first two bytes are the number of rows and the number of columns, respectively. The remaining bytes corresponds to the real pattern: 1 for a firing pixel, 0 for an off pixel. To complete the last byte, unused bits are filled with 0. The number of rows and columns is necessary to separate correctly the bit sequence according to the size of the bounding box.

proportional to this value. Sec. A.5 will show how the size of the dictionary is kept under control.

A.4 Dictionary of cluster topologies

The topology dictionary, which contains a record for each cluster topology, plays the main role in this structure. By construction, each entry of the dictionary must be uniquely identified by a key. A key must have specific characteristics. First, it must have a defined length, which must be the same for each entry of the dictionary. Second, it must be comparable with the other keys, in order to be sortable. An integer makes a good key. The raw topology is bit sequence whose length is determined by the size of the bounding box and cannot be used as a key. Fig. A.4 shows an example of a topology in its raw format.

The raw topology as a bit sequence is converted into an integer using a hash function. In this work, a C++ implementation of MurmurHash2 algorithm is used [150]. The name suggests its working principle: it is the combination of the words multiply (Mu) and rotate (r). MurmurHash2 processes the input string four bytes at a time, performing bitwise and multiplication operations, and combines all the four-byte chunks, returning an unsigned integer, i.e. a four-byte sequence. MurmurHash2 has been chosen because it provides a good uniformity of the output distribution and it is faster than other similar algorithms [151]. The uniformity of the output (or *hashcode*) distribution is an important characteristic for a hash function. Indeed, while its input can be any sequence of bits, its output is an unsigned integer. Therefore, the output can be each integer in the range $[0, 2^{32} - 1]$. If the output distribution is uniform, it is more likely that two different inputs are mapped into two different outputs.

Unfortunately, even if the selected hash algorithm has a good uniformity, it is not injective. Indeed, in rare cases ($< 1\%$), two different inputs are mapped into the same hashcode. This occurrence is called *clash*.

In this work, clashes are strongly limited by using a eight-byte extended hashcode, in which the first four bytes are the four-byte hashcode previously described and the last four bytes are the first four bytes forming the topology raw format (the first two bytes, referring to the number of rows and columns, are not considered). In this way, when two extended hashcodes are compared, a comparison of both the hashcode generated with MurmurHash2 and the first thirty-two pixels is performed. With this expedient, clashes are not mathematically avoided, but they become extremely improbable. Indeed, MurmurHash2 is characterised by an excellent avalanche behaviour, i.e. when the input changes slightly, the output changes significantly. For this reason, it is likely that two different topologies mapped into the same four-byte hashcode are extremely different. Since the second half of the extended hashcode is formed by the first thirty-two pixels of the topology, it is unlikely that two similar topologies are mapped into the same eight-byte hashcode. Extended hashcodes have been tested on 50 simulated minimum-bias Pb–Pb collisions and clashes have never been found.

In Sec. A.5, it will be explained how the size of the dictionary is kept under control.

A.5 Groups of rare topologies

As mentioned in Sec. A.3, it is necessary to limit the size of the dictionary, otherwise the compression/decompression time becomes too long. This goal is achieved through a lossy process that consists in grouping rare topologies under the same dictionary entry, where the adjective *rare* indicates topologies whose frequency is below a threshold defined *a priori*.

Groups must be formed by topologies with similar characteristics, in particular with a similar COG position and with similar uncertainties on the impact-point position. Grouping is necessarily a lossy operation, since the information about the position of the impact point of the single topology is lost in favour of average quantities. However, this loss is bearable, since it affects only rare topologies, therefore a very small fraction of the data that depends on the value of the frequency threshold.

The grouping criterion is based on the size of the bounding box. This choice was driven by the method for computing the uncertainties on the impact-point position (Eq. A.1). Indeed, topologies with a similar size have similar uncertainties on the

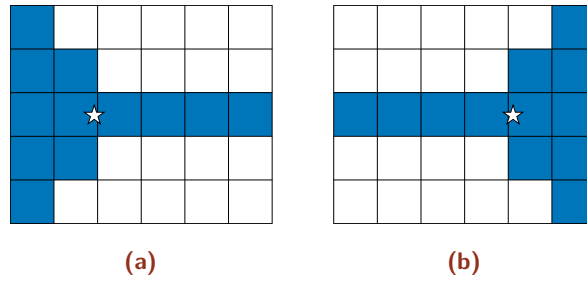


Fig. A.5.: Example of two rare topologies within the same group. The bounding boxes of (a) and (b) have the same size. However, the positions of the COG (white star) with respect to the top-left corner are very different.

impact-point position. However, even if the size is similar, the COG position of two topologies within the same group can be quite different (see Fig. A.5). This obstacle is overcome by changing the definition of the reference pixel for groups. In case of groups, the reference pixel is not the top-left corner one, but the pixel containing the COG. In this way, when storing the piece of information of a cluster with a rare topology, the position of the cluster is that of the COG, which is different for each topology of the group, and the uncertainty on the impact-point position is common for all the topologies within the group.

A.6 Structure and life of the topology dictionary

The topology dictionary, as implemented in the class `TopologyDictionary`, consists of three different elements:

1. a `std::vector` containing one entry for each common topology or group of rare topologies;
2. a `std::unordered_map` that connect an extended hash to the corresponding entry of the vector;
3. a look-up table used as shortcut for small topologies.

The first element is a vector that contains all the records for common topologies and groups of rare topologies. When storing a cluster, the topology information is expressed as the position of the corresponding element in this vector. Each element contains:

1. pattern in row format (dummy pattern for groups);
2. COG position with respect to the reference pixel (set to 0 for groups);

3. uncertainty on the position of the impact-point;
4. extended hashcode (dummy hashcode for groups);
5. number of firing pixels (dummy value for groups);
6. frequency;
7. flag for groups.

The extended hashcode, the pattern and the number of firing pixels are dummy values for groups because they are different for each topology within the same group. The vector elements are sorted in decreasing order of frequency.

If the vector is the main character of the dictionary, the other two elements play an important role in the look-up process, which is shown in flowchart in Fig. A.6. For simplicity, the look-up algorithm will be first described without considering the role of look-up table for small topologies. Referring to the flowchart in Fig. A.6, the simplified description corresponds to answering always ‘No’ to the question ‘*Does the topology fit 1 byte?*’.

The map contains an entry for each common topology. Each entry has the extended hashcode of the topology as key and the corresponding position in the vector (hereinafter simply called *ID*) as mapped value. During the look-up process, for an input topology, an extended hashcode is computed. If the topology is a common one, its extended hashcode is present in the map and the corresponding ID is found and stored. If the hashcode is not present in the table, the topology is rare and the corresponding group is computed. Each group, as well as each common topology, has its own ID and this value is stored.

The computation of the hashcode is a fast process, because it involves only arithmetical operations on a bit sequence. However, for small topologies the computation of the hashcode can be passed by using the dedicated look-up table (LUT) and the ID association can be performed with one single operation. The expression *small topology* indicates a topology that can be contained in 1 byte, i.e. with 32 pixels within the bounding box at maximum. The LUT has 8×255 elements, one for each possible topology made of 8 pixel, with a number of columns from 1 to 8. For example, in Fig. A.2 cluster C has a small topology: it has 3 columns and its raw format is 11101000. The corresponding element in the dedicated LUT is at position $(3 - 1) \times 255 = 510$. Each position of the LUT that belongs to a common topology is filled with its ID, the others are filled by default with -1. Now, it is possible to

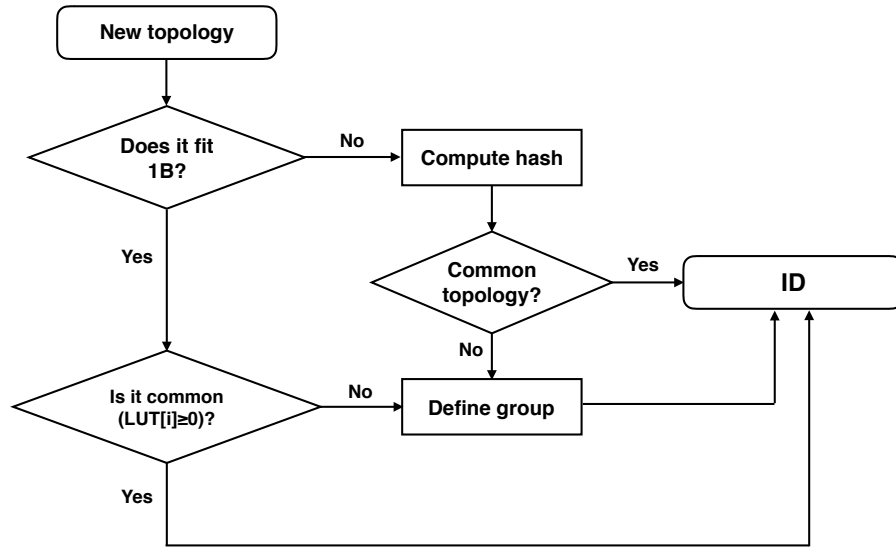
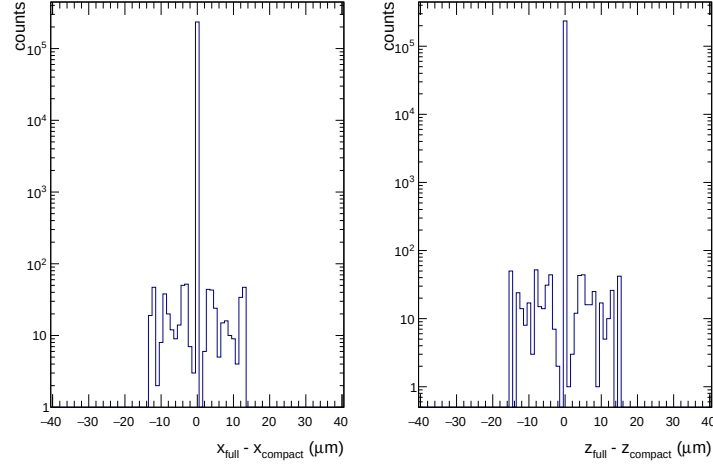


Fig. A.6.: Flowchart of the look-up algorithm.

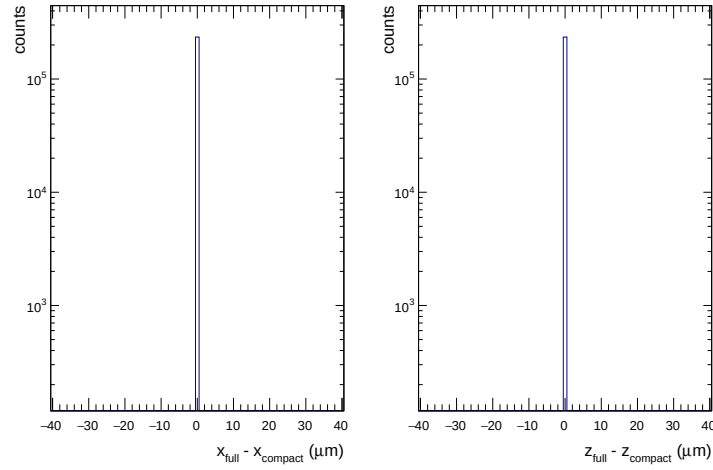
consider the whole flowchart in Fig. A.6. For an input topology, first it must be decided whether it is small or not. If it is not, the hashcode is computed and the flow is as described above. If it is a small topology, the corresponding element is searched in the dedicated LUT. If it is a common topology, the element is greater than or equal to 0 and corresponds to the ID, which is stored. Otherwise, if it is not a common topology the element is equal to -1 and the corresponding group is computed. Hence, the ID of the group is stored.

The lookup algorithm has been tested on 100 simulated minimum bias Pb–Pb collisions, corresponding to 5 million clusters. About 96% of the clusters have a small topology. The time necessary to associate to all the clusters the correct ID is 2 ms without using the LUT for small topologies and 0.7 ms using it. Therefore, the expedient of the LUT for small topologies improves the time performances by a factor of 3.

Fig. A.7 shows the difference between the cluster position obtained with the full cluster, i.e. storing the raw pattern format, and with the compact cluster, i.e. storing the topology ID. In most of the cases, there is no difference. Fig. A.7b shows that for common topologies there is no difference at all. For rare topologies, instead, a difference is present but it is always smaller than the pixel pitch. This difference is



(a) All topologies



(b) Common topologies

Fig. A.7.: Different between the COG position computed with full cluster and with compact cluster (this work) for all the topologies (a) and for common topologies only (b).

expected, because in the same group there are different topologies, for which, in general, the COG position inside the bounding box is different.

Some last words can be spent for the construction of the dictionary. This task is fulfilled by a dedicated class called `BuildTopologyDictionary`, which builds a dictionary from an input data sample of clusters.

It is important to have an updated dictionary. Indeed, the aim of this work is to compress the topology information using Huffman coding, i.e. to map the most frequent topology into the shortest bit sequence. If the data-acquisition conditions change, than the frequency distribution of topologies changes accordingly. In this case, it is necessary to create a new dictionary with the updated frequencies to

optimise the data compression. Concerning data acquisition, the idea is to have short dedicated calibration runs to update the topology dictionary and to obtain the best possible data compression.

In conclusion, the whole software for handling cluster topologies is inside O^2 , it has been tested and it works properly.

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