Universidade de São Paulo
Instituto de Física

**Medida de produção de D^0 em jatos
em colisões Pb-Pb a $\sqrt{s_{NN}} = 5.02$ TeV
com o ALICE no LHC**

Antônio Carlos Oliveira da Silva

Orientador: Prof. Dr. Alexandre Alarcon do Passo Suaide

Tese de doutorado apresentada
ao Instituto de Física para obtenção do
título de Doutor em Ciências

Banca Examinadora:

Prof. Dr. Alexandre Alarcon do Passo Suaide (IFUSP)

Prof. Dra. Ivone Freire Albuquerque (IFUSP)

Prof. Dr. Airton Deppman (IFUSP)

Prof. Dr. Eduardo de Moraes Gregores (UFABC)

Prof. Dr. David Dobrigkeit Chinellato (UNICAMP)

São Paulo

2017

FICHA CATALOGRÁFICA
Preparada pelo Serviço de Biblioteca e Informação
do Instituto de Física da Universidade de São Paulo

Silva, Antonio Carlos Oliveira da

Measurement of D^0 production in jets in Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV with ALICE at the LHC / Medida de produção de D^0 em jatos em colisões Pb-Pb a $\sqrt{s_{NN}} = 5.02$ TeV com o ALICE no LHC.
São Paulo, 2017.

Tese (Doutorado) – Universidade de São Paulo. Instituto de Física.
Depto. Física Nuclear

Orientador: Prof. Dr. Alexandre Alarcon do Passo Suaide
Área de Concentração: Física de Alta Energia

Unitermos: 1.Colisões de íons pesados relativísticos; 2.Quarks;
3.Física nuclear; 4. Partículas (Física nuclear).

USP/IF/SBI-057/2017

Universidade de São Paulo
Instituto de Física

**Measurement of D^0 production in jets
in Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV
with ALICE at the LHC**

Antônio Carlos Oliveira da Silva

Orientador: Prof. Dr. Alexandre Alarcon do Passo Suaide

Tese de doutorado apresentada
ao Instituto de Física para obtenção do
título de Doutor em Ciências

Banca Examinadora:

Prof. Dr. Alexandre Alarcon do Passo Suaide (IFUSP)

Prof. Dra. Ivone Freire Albuquerque (IFUSP)

Prof. Dr. Airton Deppman (IFUSP)

Prof. Dr. Eduardo de Moraes Gregores (UFABC)

Prof. Dr. David Dobrigkeit Chinellato (UNICAMP)

São Paulo
2017

Acknowledgements

I want to start thanking my parents, my father, Luís, who made me curious about mathematics and electronics since when I was very young, and my mother, Maria, who always made me believe that I could achieve any goal I wanted with persistence. Thank you both for loving me.

I thank my twin brother Luís, who supported me in all possible ways even before we were born. Your help made this work possible.

My supervisor, Alexandre Suaide, provided me what I call *open guidance*. He gave me freedom to work the way I thought was right and advices when I was wrong. I'm thankful for all these years that I learnt so much under his scientific supervision and for being always willing to help.

I would like to express my gratitude to all the members of the HEPIC group at the University of São Paulo: Marcelo Munhoz, Marco Bregant, Mauro Cosentino, Hugo Natal da Luz, Marcel Figueredo, Christiane Jahnke, Camila de Conti, Henrique Zanolli, Diógenes Gimenez (big adventures in Geneva), Caio Prado, Caio Laganá, Caio Eduardo, Elenos de Oliveira Filho, Douglas Vieira and Ricardo Romão. It has been a honor to be part of this group.

The whole graduation administrative staff, in special Andrea Wirkus and Éber Lima, and the nuclear physics department secretariat were always very supportive with all the administrative procedures and paperwork involved in the PhD program. I am very grateful to all of them.

During my period as PhD candidate I spent one year in Geneva, when I could work with Chiara Bianchin, who deeply introduced me to the ALICE framework and I'm very glad to have met someone so patient and skilled. There I was also lucky to live with two amazing people, Dani and Sofie, who were my landlords. Merci beaucoup!

After Geneva, I also spent almost two years in the marvelous city of Utrecht, in the Netherlands. At Utrecht university I had the pleasure of working with capable and dedicated people. Davide Lodato, Jacopo Margutti, Andrea Dubla, Alexandru Dobrin, Annelies Veen, Redmer Bertens, Sonia Vigolo, Rihan Haque, Barbara Trzeciak, Cristina Bedda, Jasper van der Maarel, Alberto Caliva, Emilia Leogrande, Darius Keijder (Bedankt voor het helpen met Nederlands), Lennart van Doremalen, Sandro Bjelogrić, Luuk Vermunt, Syaefudin Jaelani, Tuva Herenui Richert Shaw, Chunhui Zhang, Mike Sas, Naghmeh Mohammadi, Hongkai Wang, Alessandro Grelli, Thomas Peitzmann,

Ton van den Brink, Marco van Leeuwen, Panos Christakoglou, Paul Kuijer, Raimond Snellings and, in special, to my supervisor for the time I was abroad, André Mischke, who gave me support and guidance.

This work would not be possible without the support of Heavy-Flavour Correlations and Jets working group, in special, I want to thank Elena Bruna, Andrea Rossi, Fabio Collamaria, Salvatore Aiola and, again, Barbara Trzeciak.

I also have friends I would like to thank. Henrique Goes, Henrique Machado, Maurício Borges, Artur Fontainha, Kathe Colomba, Isadora Leite, Karina Cecchinell, Maya Prisyazhnaya, Giuliana Ratzel, Lucia Sestokas and my cousin Luiz Dario. Life is better with all of them.

I thank also Rodrigo Uchida Ichikawa and Camila Machado for the friendship and for the discussions about science and life.

I thank Aline Oliveira, who makes me happy. Her support and kindness allowed me to finish this work. Thank you for the feelings, companionship and love.

I thank FAPESP for the financial support of this work.

*“Tho’ much is taken, much abides; and tho’
We are not now that strength which in old days
Moved earth and heaven, that which we are, we are;
One equal temper of heroic hearts,
Made weak by time and fate, but strong in will
To strive, to seek, to find, and not to yield.”*
— Ulysses by Alfred, Lord Tennyson

Abstract

Charm quarks are created in the early stages of heavy-ion collisions in hard-scattering processes. Therefore, they are ideal probes of the Quark-Gluon Plasma (QGP), which is a state of matter where partons contained in hadrons, at high temperatures (~ 150 MeV) or density (about five times the density of ordinary matter), change to a deconfined state of quarks and gluons. The fragmentation of charm quarks can produce D mesons.

Jets containing a D meson as one of their constituents can be identified as originating from heavy-quark fragmentation. These jets are a valuable tool to characterize the charm interaction with the QGP. Charmed jets can provide information to the study the mass-dependent energy loss by the measurement of the modification of the charm-jet yield in Pb-Pb collisions with respect to pp collisions as a function of the jet transverse momentum. Moreover, a further insight can be obtained with the measurement of the momentum-fraction distribution, which is of particular interest to investigate the possible influence of the medium in the charm-jet fragmentation.

D mesons are reconstructed through an invariant mass analysis of their hadronic decay channels, rejecting the large combinatorial background with topological selections exploiting the relatively large lifetime of D mesons and the particle identification capabilities of the ALICE detector. Jets are reconstructed with the anti- k_T algorithm using D-meson candidates and charged tracks.

The measurement of the transverse momentum spectrum of jets containing D mesons in Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV will be presented. These results lead to new possibilities of probing the QGP properties when compared to baseline measurements. Furthermore, the methods developed can be employed in the study of charm-quark fragmentation functions in heavy-ion collisions.

Resumo

Quarks charm são criados em estágios iniciais da colisão de íons pesados em processos de espalhamento duro. Portanto eles são sondas ideais para o Plasma de Quarks e Glúons (QGP), que é um estado da matéria em que os partons contidos em hadrons, em condições de alta temperatura (~ 150 MeV) ou densidade (cerca de cinco vezes a densidade da matéria ordinária), passam para um estado livre de quarks e glúons. A fragmentação de quarks charm pode produzir mésons D.

Jatos contendo um méson D como um de seus constituintes podem ser identificados como originados de uma fragmentação de quark pesado. Jatos contendo mésons D são valiosas ferramentas para caracterizar interações de quarks charm com o plasma. Jatos de charm podem fornecer informações para o estudo da perda de energia dependente da massa pela medida da modificação da produção de jatos de charm em colisões de núcleos de chumbo com respeito a colisões entre prótons em função do momento transversal dos jatos. Além disso, uma visão mais profunda pode ser obtida com a medida da distribuição da fração de momento, que é de particular interesse para investigar a possível influência do meio na fragmentação de quarks charm em jatos.

Mésons D são reconstruídos através da análise de massa invariante de seu canal de decaimento hadrônico, rejeitando uma grande quantidade de fundo combinatório com seleções topológicas e explorando o tempo de vida relativamente longo de mésons D e as capacidades de identificação de partículas do detector ALICE. Jatos são reconstruídos com o algoritmo anti- k_T usando candidatos a mésons D e partículas carregadas.

A medida do espectro de momento transversal de jatos contendo mésons D em colisões de núcleos de chumbo a energias de $\sqrt{s_{NN}} = 5.02$ TeV será apresentada. Esses resultados levam a novas possibilidades de sondagem das propriedades do QGP quando comparados com medidas de referência. Por fim, os métodos desenvolvidos podem ser empregados no estudo das funções de fragmentação de quarks charm em colisões de íons pesados.

Contents

List of Figures	xiii
List of Tables	xvii
1 Introduction	1
1.1 The Internal Structure of Nucleons	1
1.2 Gluons	2
1.3 Standard Model	4
1.4 Quantum Chromodynamics	5
1.5 Quark-Gluon Plasma	8
2 Signatures of the Quark-Gluon Plasma	15
2.1 Geometry of Heavy-Ion Collisions	15
2.2 Observables in Heavy-Ion Collisions	16
2.3 Collective Effects	17
2.4 Energy Loss	18
2.4.1 Collisional Energy Loss	18
2.4.2 Radiative Energy Loss	20
2.4.3 Nuclear Modification Factor	24
2.5 Quark Fragmentation	25
2.6 Jets	27
3 Heavy Flavour	29
3.1 Production Mechanisms	29
3.2 Heavy-Flavour in Heavy-Ion Collisions	32
4 ALICE Experiment	35
4.1 Large Hadron Collider	35
4.2 ALICE	36
4.2.1 Inner-Tracking System	38
4.2.2 Time-Projection Chamber	39
4.2.3 Time-of-Flight	40
4.2.4 VZERO	41

5	Analysis	43
5.1	Event Selection	43
5.2	D-Meson Reconstruction	44
5.3	Jet Reconstruction	46
5.3.1	Background Subtraction	47
5.4	D-Tagging	49
5.5	Signal Extraction	51
5.5.1	D-Tagged Jet p_T Direct Extraction	51
5.5.2	Side-Band Subtraction	51
5.5.3	Method Comparison	53
5.6	Efficiencies Corrections	53
5.6.1	D-Meson Reconstruction Efficiency	55
5.7	Unfolding	57
5.7.1	Bayesian Method	58
5.7.2	SVD Method	59
5.7.3	Challenges Concerning Unfolding	60
5.7.4	Response Matrix	64
5.7.5	Unfolding Closure Tests	66
5.8	Feed-Down	68
5.8.1	Unfolding Results	72
5.9	Systematic Uncertainties	74
5.9.1	Multi-Trial Signal Extraction	74
5.9.2	Feed-Down	76
5.9.3	Unfolding	78
5.9.4	Final Systematic Uncertainties	79
6	Conclusions	81
6.1	The Measurement	81
6.2	Method Development	83
6.3	Final Considerations	85
A	Conventions and Notations	87
	Bibliography	89

List of Figures

1.1	Quark-Antiquark production in e^+e^- annihilation.	2
1.2	Jet angles and Ellis-Kaliner angle $\tilde{\theta}$. $\tilde{x}_{1,2,3}$ are the energy fractions of the jets in the Lorentz transformed frame.	3
1.3	Distribution of $\cos\tilde{\theta}$. The solid and dotted lines are the QCD calculation for vector and scalar gluons respectively.	4
1.4	Constituents of matter according to the Standard Model.	5
1.5	Compilation of data on the measurements of $\mathbf{R}$	6
1.6	Equation of state for a pion gas (blue) and a gas of quarks and gluons (red). The crossing point determines the critical temperature of transition.	11
1.7	Quark-Gluon Plasma phase diagram obtained by equation 1.43.	13
2.1	Schematic representation of a heavy-ion collision.	16
2.2	Schematic representation of the QGP space-time evolution.	17
2.3	Schematics of a non-central heavy-ion collision (left) and anisotropic momentum distribution due to pressure gradients (right).	17
2.4	Measurement of ν_2 for charged particles (blue) and the average of D mesons species (black) measured by the ALICE Collaboration.	18
2.5	Feynman diagrams of one gluon radiation.	22
2.6	Radiation amplitude as a function of θ for <i>down</i> (continuous line) and <i>charm</i> quark (dashed line). The parameters for the calculation are $p_i = (0, 0, 10)$ GeV, $p_f = (0.3, 0.2)$ GeV, $\mu_D = 0.5$ GeV for the quark and $\omega = 0.001$ GeV, $\phi = \pi/2$ for the radiated gluon.	23
2.7	Nuclear modification factor of D mesons (average of D^0 , D^+ and D^{*+}), charged pions and charged particles in 0-10% central Pb-Pb events at $\sqrt{s_{NN}} = 2.76$ TeV.	24
2.8	Jet pictured as particles inside a cone.	27
2.9	Jet nuclear modification factor in 0-10% (left) and 10-30% (right) centrality of Pb-Pb collisions at $\sqrt{s} = 2.76$ TeV with respect to the same measurement done in pp collisions. Data is compared to JEWEL (black line) and YaJEM (green dashed line) models.	28
3.1	Flavour Creation diagrams. Two heavy quarks in the hard subprocess.	30
3.2	Flavour Excitation diagrams. One heavy quark in the hard subprocess.	30
3.3	Gluon Splitting diagrams. No heavy quarks in the hard subprocess.	31

3.4	Charm (a) and bottom (b) production cross sections in pp collisions.	32
3.5	Average R_{pPb} of D mesons in pPb collisions at $\sqrt{s_{NN}} = 5.02$ TeV compared with models. Statistical (bars) and systematic (boxes) uncertainties are also presented. . .	33
4.1	Schematic representation of the Large Hadron Collider complex.	36
4.2	Picture of the ALICE detector and subsystems.	37
4.3	The six layers of the Inner-Tracking System.	38
4.4	The six layers of the Inner-Tracking System.	39
4.5	Average energy loss per unit of path length (arbitrary units) as a function of the transverse momentum over the particle charge.	40
4.6	TOF β in function of particle momentum.	41
4.7	V0A and V0C schematic rings.	42
4.8	Number of events as a function of the V0 amplitude. The shaded areas mark regions attributed to centrality percentiles. The fit (in red) is performed using Glauber model. . .	42
5.1	Number of events as a function of the event centrality determined by V0 detector. . .	44
5.2	Schematics of the D^0 decay by the hadronic channel.	46
5.3	Inclusive jet transverse momentum distribution from the 0-20% most central Pb-Pb events at $\sqrt{s_{NN}} = 5.02$ TeV. The distribution is background subtracted.	48
5.4	The δp_T distributions provides the magnitude of the jet background fluctuation estimated with random cones method.	50
5.5	Representation of the method applied to reconstruct heavy-flavour jets. The daughter tracks are replaced by the D^0 candidate itself.	50
5.6	Invariant mass distributions in intervals of jet p_T from -10 to 35 GeV/c. D-meson transverse momenta are in the range of 3 to 20 GeV/c.	52
5.7	Invariant mass distributions in intervals of jet p_T	53
5.8	Jet transverse momentum spectra for different intervals of D^0 p_T . The red points corresponds to the signal region and the blue points to the normalized side-band region.	54
5.9	Jet transverse momentum spectra for different intervals of D^0 p_T	54
5.10	D-tagged jets transverse momentum. Jet p_T direct extraction method are the red points and the blue is the distribution using the side-band method.	55
5.11	Prompt D meson reconstruction efficiency in Pb-Pb collisions at $\sqrt{s} = 5.02$ TeV. The statistical uncertainties are smaller than the points.	56
5.12	D-tagged jet transverse momentum for both methods. The red distribution was obtained using direct extraction. The side-band method is pictured with blue points. .	57
5.13	Relative statistical uncertainties of D-tagged jet p_T for the two methods.	58
5.14	Signal jet with constituents in red is presented on the left. On the right, the jet is split into two jets due to the presence of soft uncorrelated particles pictured in blue. .	61
5.15	Tracks from simulation are in red and tracks from data are in blue.	61
5.16	Number of tracks ratio between PYTHIA tracks from the embedded and simulated-only samples.	62

5.17	Transverse momentum fraction between PYTHIA tracks from the embedded and simulated-only samples.	63
5.18	Inclusive jet transversal momentum distributions from Pb-Pb data in blue and embedding (PYTHIA pp simulations and Pb-Pb data) in red.	64
5.19	Prompt D^0 reconstruction efficiency in pp collisions at $\sqrt{s} = 5.02$ TeV.	65
5.20	Response matrix for prompt D-tagged jet unfolding.	66
5.21	Left: true (red line) and unfolded (blue points) D-tagged jet transversal momentum distributions. Right: relative difference between true and unfolded spectra.	66
5.22	Left: relative difference between true and unfolded spectra. The z axis is the number of entries. Right: standard deviation of the relative difference distribution. The unfolding procedure was performed using 500 spectra.	67
5.23	D^0 from B decays reconstruction efficiency in Pb-Pb collisions at $\sqrt{s} = 5.02$ TeV. . .	68
5.24	D^0 from B decays reconstruction efficiency in pp collisions at $\sqrt{s} = 5.02$ TeV. . . .	69
5.25	Response matrix built using only D^0 from B mesons.	70
5.26	Fraction of non-prompt over measured D-tagged jets as a function of jet p_T	71
5.27	Measured prompt D-tagged jet p_T distribution corrected by D-meson reconstruction efficiency and feed-down from B meson decay. The entries are divided by the bin width. . .	71
5.28	Unfolded prompt D-tagged jet p_T spectrum.	72
5.29	Refolded (blue line) and measured (blue points) distributions.	73
5.30	Relative systematic uncertainties of the D-tagged jet signal extraction using the direct extraction and side-band methods.	75
5.31	Non-prompt D-tagged jet relative difference between the variation of the POWHEG simulation over the central points as a function of jet p_T	76
5.32	Relative systematic uncertainties of the non-prompt D-meson R_{AA} hypothesis. . . .	77
5.33	Relative systematic uncertainties from Unfolding.	78
5.34	Final combined relative systematic uncertainties.	79
6.1	Invariant yield of D^0 -tagged jet transverse momentum spectrum. The red points are data and the red boxes are the systematic uncertainties. The simulation is presented in blue points with uncertainties represented by blue boxes.	82
6.2	Nuclear modification factor of D-tagged jets in function of the jet transverse momentum. Red points are the measurement over the simulated baseline. The red boxes are the systematic uncertainties from the measurement and the blue boxes represent systematic uncertainties from the simulation.	83
6.3	Comparison between data and simulation with D-meson R_{AA} . The red points are data and the red boxes are the systematic uncertainties. The simulation is presented in blue points with uncertainties represented by blue boxes.	84
A.1	Geometry convention in ALICE experiment.	87

List of Tables

4.1	Dimensions and two-track resolution of the ITS layers.	38
5.1	Kinematical and topological cuts for D^0 candidates in Pb-Pb collisions at $\sqrt{s_{\text{NN}}} =$ 5.02 TeV.	45

Chapter 1

Introduction

“There is no knowledge that is not power.”

— Ralph Waldo Emerson

Since a long time, mankind makes questions about the origin of the things. The *Big-Bang* is today the most acceptable theory of how the universe began and, consequently, matter was created. Science was able to provide a considerable amount of evidence like Cosmic Microwave Background Radiation (radiation from the epoch of recombination and subsequently photon decoupling), Large-Scale Structure (mass distribution in the universe) and Hubble’s Law (expanding universe) that corroborate this description of our cosmos starting as a hot and dense explosion.

1.1 The Internal Structure of Nucleons

The discovery of the electron and the nucleus last century was a significant advance in the understanding of the matter. Later, quantum mechanics changed again the way science described atoms and subatomic particles introducing new ideas such as particle-wave duality, the uncertainty principle and antimatter.

Between the decades of 1950s and 1960s, many new subatomic particles were discovered due to the use of particle accelerators and a theory that could explain the existence of such number of different entities was required at that time. Gell-Mann proposed in 1964 that all particles observed so far could be explained by the combination of three elementary particles of spin 1/2 with electric charge equals to 2/3 and -1/3 of the positron. These particles were called *quarks* by the proposer of the model. Even so, it was necessary to introduce a new quantum number, later known as colour.

In 1969, deep inelastic scattering experiments using electrons to probe protons at the Stanford Linear Accelerator Center (SLAC) observed a scaling behavior consistent with the *Bjorken scaling* [1, 2], in which, if $Q^2 \gg M^2$, the structure functions of the proton can be approximated as

$$F(x, Q^2) \approx F(x), \quad (1.1)$$

i.e., the structure function is nearly independent of the square of the transferred momentum-energy

$$Q^2 = -q^2 = -(k - k')^2, \quad (1.2)$$

where k is the initial and k' the final 4-momentum of the electron. The variable x is called Bjorken- x

and is defined as

$$x \equiv Q^2/2p \cdot q, \quad (1.3)$$

with p the 4-momentum of the proton.

The data collected by SLAC could be explained with the existence of an internal structure of the nucleon, i.e., the existence of three scattering centers. These constituents of protons and neutrons were called partons.

Though it was known that partons were $1/2$ -spin particles, their association with Gell-Mann's quarks was not direct. It came only later with the confirmation of the squared charges of the partons as $1/9e^2$ and $4/9e^2$ by other experiments [3, 4, 5, 6].

1.2 Gluons

In the late 1970s, three-jet events were observed in electron-positron collisions at PETRA by the TASSO [7], MARK J [8], PLUTO [9] and JADE [10] experiments. Figure 1.1a pictures the production of a quark-antiquark pair in an electron-positron annihilation. The quark fragmentation originates jets of particles produced back-to-back.

Figure 1.1: Quark-Antiquark production in e^+e^- annihilation.

There is a probability of the quark (or antiquark) to radiate a high-energy gluon before the fragmentation with an angle large enough to be identified as an independent jet, as pictured in figure 1.1b.

Due to the gluon emission in three-jet events, the jet energy fraction with respect to the beam energy ($x_i = E_i/E_b$) can be ordered as $x_1 \geq x_2 \geq x_3$. Considering the approximation where quark masses are negligible, x_i can be written as

$$x_i = \frac{E_i}{E_b} = \frac{2 \sin \theta_i}{\sin \theta_1 + \sin \theta_2 + \sin \theta_3} \quad (1.4)$$

where θ_1 , θ_2 and θ_3 are defined in figure 1.2a and E_b is the energy of the electron (or positron) beam.

Figure 1.2: Jet angles and Ellis-Kaliner angle $\tilde{\theta}$. $\tilde{x}_{1,2,3}$ are the energy fractions of the jets in the Lorentz transformed frame.

A Lorentz transformation from the laboratory rest frame to the center-of-mass of jets 2 and 3 is done boosting the event in the direction of jet 1. Jets 2 and 3 are then back-to-back and the Ellis-Kaliner angle $\tilde{\theta}$ is defined as the angle that these jets make with the boost axis as shown in figure 1.2b.

Neglecting again the quark masses, the $\cos \tilde{\theta}$ can be written as

$$\cos \tilde{\theta} = \frac{\sin \theta_2 - \sin \theta_3}{\sin \theta_1}. \quad (1.5)$$

The angular distribution of the jets as explained above depends on the gluon spin [11]. The results from the TASSO collaboration [12] are shown in figure 1.3. The data is compared to the QCD calculation for scalar and vector gluons and clearly favors spin-1 gluons.

The discovery of the gluon was a decisive test of the *Standard Model* (SM), which is the theory of the elementary particles and their interactions. Gluons are the carriers of the strong force described by *Quantum Chromodynamics* (QCD) [13, 14, 15].

Figure 1.3: Distribution of $\cos \tilde{\theta}$. The solid and dotted lines are the QCD calculation for vector and scalar gluons respectively.

1.3 Standard Model

The Standard Model is a gauge quantum field theory with internal symmetry $SU(3) \times SU(2) \times U(1)$ (Yang-Mills theory) that describes the elementary constituents of matter (fermions) and their fundamental interactions mediated by gauge bosons.

The particles that compose the SM are summarized in figure 1.4 [16]. Each fermion has its respective antiparticle with electric and colour charge with opposite sign.

The $SU(2) \times U(1)$ gauge group is the internal symmetry of the electroweak interaction. Spontaneous symmetry breaking of the electroweak symmetry leads to the W^+ , W^- and Z^0 bosons of the weak force and the photon (γ), the gauge boson of the electromagnetic interaction. The electroweak theory, also known as Glashow-Weinberg-Salam model, succeeded in unifying two fundamental forces.

The theory of strong interactions contains the $SU(3)$ symmetry and will be discussed in section 1.4. The SM gives a good description of three of the four known fundamental forces, saying nothing about the gravitational interaction.

The last ingredient of the Standard Model is the Higgs boson. The existence of this particle was only recently confirmed by CERN [17]. The Higgs boson is a quantum of the Higgs field that provides mass to fermions and the bosons of the weak force via the Higgs Mechanism (though the mass of neutrinos is a more complicated case).

Figure 1.4: Constituents of matter according to the Standard Model.

1.4 Quantum Chromodynamics

Quantum Chromodynamics (QCD) is a non-abelian (non-commutative) quantum field theory with $SU(3)$ symmetry. This symmetry introduces three quantum numbers called colour charges red (R), green (G) and blue (B). Antiquarks will have colour with negative sign, or antired ($\bar{R}$), antigreen ($\bar{G}$) and antiblue ($\bar{B}$) anticolours.

The model proposed by Gell-Mann could explain hadrons with the introduction of a colour quantum number for quarks. In order to form a bound state, the combination of quarks has to be colourless. Baryons are composed of three quarks with different colours and the combination of the three colours results in white. For mesons, since they are formed by a quark and an antiquark, if the quark has a given colour, the antiquark has to be of the same colour with opposite sign.

Quarks have never been found free. Hadrons are observed only in colour singlet states, which is known as *colour confinement*. There is no analytical or experimental proof of confinement, though observations are compatible with this idea.

The number of colours was confirmed experimentally by measuring the ratio of the cross sections of hadronic over muonic productions in e^+e^- annihilations as shown in equation 1.6.

$$R = \frac{\sigma(e^+e^- \rightarrow H)}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} = \left(\sum_k e_q^2 \right) N_C \quad (1.6)$$

where N_C is the number of colours, k is the quark flavour and e_q is the quark electric charge. This ratio is proportional to the number of colours [18] neglecting finite quark mass and QCD corrections.

The ratio R depends on the collision energy in the center-of-mass system reference frame ($\sqrt{s}$).

Inside the interval $1.3 \lesssim \sqrt{s} \lesssim 3$ GeV only quarks up, down and strange are expected to be produced so, considering $N_C = 3$, a value of $R = 3(4/9 + 1/9 + 1/9) = 2$ should be observed. Figure 1.5 contains a compilation of measurements of R [19] and is in good agreement with the existence of 3 colour charges. The same can be seen for the regions where charm and bottom quarks are produced.

Figure 1.5: *Compilation of data on the measurements of R .*

All the properties of the strong force mentioned so far have to be accommodated by Quantum Chromodynamics. The Lagrangian of the theory is written as

$$\mathcal{L}_{QCD} = -\frac{1}{2} \text{tr} (G_{\mu\nu} G^{\mu\nu}) + \sum_k^{n_f} \bar{\psi}_k^i (i\gamma^\mu D_\mu)^{ij} \psi_k^j - \bar{\psi}_k^i m_k \delta^{ij} \psi_k^j \quad (1.7)$$

where ψ_k^i is a k -flavoured quark ($k = u, d, s, \dots$) with colour i and

$$G_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu - ig [A_\mu, A_\nu], \quad (1.8)$$

$$D_\mu \psi_k = (\partial_\mu - ig A_\mu) \psi_k \quad \text{and} \quad (1.9)$$

$$A_\mu = \sum_{a=1}^8 A_\mu^a \frac{\lambda^a}{2}. \quad (1.10)$$

The λ^a are the matrices that satisfy the SU(3) commutation relations

$$\left[\frac{\lambda^a}{2}, \frac{\lambda^b}{2} \right] = if^{abc} \frac{\lambda^c}{2} \quad (1.11)$$

with the normalization condition

$$\text{tr}(\lambda^a \lambda^b) = 2\delta^{ab} \quad (1.12)$$

where f^{abc} is the antisymmetric structure constant of SU(3) and δ^{ab} is the Kronecker delta. The three-by-three (in colour space) matrix A_μ is composed of the gluon fields A_μ^a and the generators of the SU(3) gauge group ($\lambda^a/2$). The index a is a combination of colour and anticolour.

Since there are 8 generators in the SU(3) group, QCD provides a description of strong interactions with 8 gluons. The non-commutativity of the theory (equation 1.11) successfully describes some features of the strong force, e.g., gluons carry colour and anticolour charge and therefore interact with themselves. From these self-interacting gluons arises *asymptotic freedom*.

In equations 1.8 and 1.9, the constant g is the gauge coupling related to the coupling constant of the strong force α_s (equation 1.13), which is a parameter that gives the strength of the strong interaction.

$$g^2 = 4\pi\alpha_s \quad (1.13)$$

The renormalization group β -function in a non-abelian gauge theory in a one-loop level can be written as

$$\beta(g) = -\frac{g^3}{(4\pi)^2} \left(\frac{11}{3}C_2(G) - \frac{4}{3}n_f C(r) \right) \quad (1.14)$$

where n_f is the number of flavours, $C_2(G)$ in a SU(N) group is equal to N ($N_C = 3$ for QCD) and $C(r)$ is defined such that

$$\text{tr}(t_a t_b) = C(r)\delta_{ab} \quad (1.15)$$

with $t_a = \lambda_a/2$. From equations 1.12 and 1.15 one concludes that $C(r) = 1/2$. So the β -function for the QCD is written as

$$\beta(g) = -\frac{g^3}{(4\pi)^2} \left(11 - \frac{2}{3}n_f \right) = -bg^3. \quad (1.16)$$

Since there are six known flavours of quarks, in the context of QCD $b > 0$. This defines the behavior of the gauge coupling in the evolution of the momentum scale of the strong interaction $q \rightarrow \lambda q$. The gauge coupling has to obey the equation

$$\frac{dg(\rho)}{d\rho} = -bg(\rho)^3 \quad (1.17)$$

where $\rho = \ln \lambda$. Solving equation 1.17 leads to

$$g^2(\rho) = \frac{g_0^2}{1 + 2bg_0^2\rho}. \quad (1.18)$$

with $g_0 = g(0)$. Analysing the denominator of 1.18 it is straightforward to note that $g(\rho) \rightarrow 0$ when $\lambda \rightarrow \infty$. This means that for large momentum scale the coupling constant vanishes logarithmically leading to asymptotic freedom.

It is convenient to write λ as a function of the ratio of the momentum involved in the interaction Q and the renormalization scale μ_R as $\lambda^2 = Q^2/\mu_R^2$ and, using 1.13, the expression for α_s becomes

$$\alpha_s(Q^2) = \frac{\alpha_s(\mu_R^2)}{1 + 4\pi b\alpha_s(\mu_R^2) \ln Q^2/\mu_R^2}. \quad (1.19)$$

Equation 1.19 can be simplified by defining a parameter Λ such that

$$\ln \Lambda^2 = \ln \mu_R^2 - \frac{1}{4\pi b\alpha_s(\mu_R^2)} \quad (1.20)$$

and α_s can be written as

$$\alpha_s(Q^2) = \frac{4\pi}{(11 - \frac{2}{3}n_f) \ln Q^2/\Lambda^2}. \quad (1.21)$$

The parameter Λ is called *QCD scale constant*. For $Q^2 \gg \Lambda^2$ the coupling constant becomes small and QCD can be treated perturbatively. However, for $Q^2 = \Lambda^2$, the coupling constant diverges. Nevertheless, the expression for α_s was obtained from the one-loop approximation, so it is expected that the coupling constant diverges in non-perturbative regimes.

It is worth noting that $\alpha_s(\mu_R^2)$ is defined by Λ for a given approximation of n-loops. Also, the value of Λ depends on how the renormalization group equation is solved. For these reasons it is very common that the coupling constant α_s is presented for a given Q^2 value which is chosen to be the mass of the Z^0 boson (M_Z). Using the modified minimal subtraction ($\overline{MS}$) the QCD constant is commonly quoted as $\Lambda = 200$ MeV.

From equation 1.21 it is clearly seen that for short distances ($Q^2 \gg \Lambda^2$) the strong interaction becomes small and asymptotic freedom arises. This is the opposite screening effect found in *Quantum Electrodynamics* (QED) called antiscreening. The effect of quark loops (screening) and gluon loops (antiscreening) is dominated by the latter as long as the number of flavours is smaller than 17.

In light of the properties of strong interactions and, specially, asymptotic freedom, it is expected that, for a sufficiently high temperature, hadrons should change to a state where quarks are so weakly interacting that they become free and the system presents partonic degrees of freedom. This state of matter is called the *Quark-Gluon Plasma* (QGP).

1.5 Quark-Gluon Plasma

Quantum Chromodynamics predicts that quarks, under certain conditions, are weakly interacting and no longer confined in hadrons. This state of matter is known as *Quark-Gluon Plasma* (QGP). *Lattice QCD*[20] (lQCD) calculations also predict a critical temperature of transition between deconfined strong interacting matter and a hadronic state.

A deconfined state of matter should present additional partonic degrees of freedom (colour, flavour, etc) that will increase the entropy of the system. Using a very simple model [21] it is possible to predict the critical temperature of transition. This model compares a relativistic pion gas pressure and a relativistic gas of deconfined quarks and gluons.

Thermodynamics shows that systems choose the state of minimum free energy, which means maximum entropy. The temperature of transition can be found where the quark-gluon pressure becomes equal to the pressure of hadrons.

Considering a relativistic gas of particles in a volume V , the thermodynamical potential in the grand canonical ensemble is given by equation 1.22.

$$\phi = -\frac{1}{\beta} \int_0^\infty \ln [1 \pm f(\epsilon)] \mathcal{D}(\epsilon) d\epsilon \quad (1.22)$$

where $\mathcal{D}(\epsilon)$ is the density of states

$$\mathcal{D}(\epsilon) = \frac{Vd}{(2\pi)^3} \epsilon^2, \quad (1.23)$$

in which d is the number of degrees of freedom (or degeneracy).

The average number of particles in a given energy ϵ is defined as

$$f(\epsilon) = \frac{1}{e^{\beta(\epsilon-\mu)} \pm 1} \quad (1.24)$$

where $1/\beta = T$ is the temperature of the system. In both equations 1.23 and 1.24 the $\pm$ symbol indicates that particles can be bosons (positive sign) and follow the Bose-Einstein distribution or fermions (negative sign) and obey the Fermi-Dirac distribution.

Pressure and internal energy density of the systems are connected and this is conveniently represented by the grand canonical potential. Considering, at first, bosons in the simplest case where $\mu = 0$, the internal energy is

$$U = \int_0^\infty f(\epsilon) \mathcal{D}(\epsilon) \epsilon d\epsilon, \quad (1.25)$$

and integrating the grand canonical potential by parts it becomes:

$$\phi = \frac{1}{\beta} \frac{4\pi d}{3(2\pi)^3} \ln [1 + f(\epsilon)] \epsilon^3 \Big|_0^\infty - \frac{4\pi d}{3(2\pi)^3} \int_0^\infty f(\epsilon) \epsilon^3 d\epsilon = -\frac{1}{3} \int_0^\infty f(\epsilon) \mathcal{D}(\epsilon) \epsilon d\epsilon \quad (1.26)$$

The surface term vanishes and the integral is clearly related to the internal energy by $U = -\frac{1}{3}\phi$. Since the pressure also has a simple expression given by $\phi = -pV$, it can be written as:

$$p = -\frac{\phi}{V} = \frac{1}{3} \frac{U}{V} = \frac{1}{3} u \quad (1.27)$$

The internal energy density u for a gas of hadrons can be calculated in first approximation as a gas of pions

$$u = \frac{d}{2\pi^2} \int_0^\infty \frac{\epsilon^3}{e^{\beta\epsilon} - 1} d\epsilon = 3 \times \frac{\pi^2}{30} T^4. \quad (1.28)$$

The number of degrees of freedom for the pion gas is 3 due to the three types of pions (π^+ , π^- and π^0). From equation 1.27 the pressure of a gas of pions is

$$p_H = \frac{\pi^2}{30} T^4. \quad (1.29)$$

The internal energy of the quark-gluon plasma contains a contribution from the quarks and another from gluons. Since gluons are also bosons it is straightforward to calculate the internal energy

$$u_g = 8_{color} \times 2_{spin} \times \frac{\pi^2}{30} T^4 \quad (1.30)$$

and pressure

$$p_g = \frac{8\pi^2}{45} T^4 \quad (1.31)$$

for gluons taking into account that gluons have 8 degrees of freedom of color charge and 2 of spin.

The internal energy density and pressure of the QGP also has a quark contribution and, since they are fermions, it is easy to show that its internal energy density differs from that of bosons from a factor 7/8 (not taking into account the degeneracy) as seen in equation 1.32.

$$\frac{\int_0^\infty \frac{\epsilon^3}{e^{\beta\epsilon}+1}}{\int_0^\infty \frac{\epsilon^3}{e^{\beta\epsilon}-1}} = \frac{\frac{7\pi^4}{120}}{\frac{\pi^4}{15}} = \frac{7}{8} \quad (1.32)$$

Using the result above the quark-gluon internal energy density is given by equation 1.33.

$$u_{QGP} = \left\{ \frac{7}{8} [2_{q\bar{q}} \times 2_{spin} \times 2_{flavour} \times 3_{color}] + 2_{spin} \times 8_{color} \right\} \frac{\pi^2}{30} T^4 \quad (1.33)$$

The degrees of freedom of quarks are 2 for the quark-antiquark, 2 for the two half-integer possible spins, 2 for the flavours (considering only u and d quarks) and 3 for the three different color charges. There is still one last term that must be added to the QGP internal energy density, which is the bag pressure B . It is related to the difference of energy density between the normal QCD-vacuum outside and the perturbative QCD-vacuum inside the hadron ($\epsilon_{vacuum}^{in} - \epsilon_{vacuum}^{out} \equiv B$).

The bag pressure can be calculated by minimizing the energy of a system of N confined quarks in a spherical volume of radius R [22]

$$E = \frac{2.02N}{R} + \frac{4\pi}{3} R^3 B. \quad (1.34)$$

Applying the condition $dE/dR=0$ the bag pressure can be written as

$$B^{1/4} = \left(\frac{2.04N}{4\pi} \right)^{1/4} \frac{1}{R}. \quad (1.35)$$

With this last ingredient we can write the QGP pressure using 1.27 and 1.33 as:

$$p_{QGP} = 37 \frac{\pi^2}{90} T^4 - B. \quad (1.36)$$

As said before, the critical temperature T_C can be found invoking the condition $p_{qg} = p_\pi$. Using equations 1.29 and 1.36 we write an expression for the temperature of transition (eq. 1.37):

$$37 \frac{\pi^2}{90} T_C^4 - B = 3 \frac{\pi^2}{90} T_C^4 \Rightarrow T_C \simeq 0.72 B^{1/4}. \quad (1.37)$$

Considering 3 quarks inside a bag of radius $R=0.8$ fm the bag pressure is

$$B^{1/4} = 206 \text{ MeV}, \quad (1.38)$$

which results in a critical temperature of

$$T_C = 144 \text{ MeV.} \quad (1.39)$$

At this temperature the partons are no longer confined in hadrons. This is remarkably close to recent results from *lQCD* calculations [23]. The condition used in equation 1.37 is pictured in figure 1.6 using equations 1.29 and 1.36.

Figure 1.6: Equation of state for a pion gas (blue) and a gas of quarks and gluons (red). The crossing point determines the critical temperature of transition.

Considering now the case where the chemical potential is not zero, it is possible to write a closed form for the Helmholtz free energy density considering gluons and two flavours ($k=u,d$) of massless quarks and antiquarks [24] as

$$f = -\frac{8}{45}\pi^2 T^4 + \sum_k \left(\frac{3}{4\pi^2} \mu_k^4 + \frac{1}{2} \mu_k^2 T^2 - \frac{7}{60} \pi^2 T^4 \right). \quad (1.40)$$

The condition for the presence of the QGP is slightly different from the previous case. It is now considered that the QGP pressure has to be equal or higher than the QCD vacuum pressure

$$p_{QGP} = B \quad (1.41)$$

Using the thermodynamical relations $f = u - Ts$ and $s = -\partial f / \partial T$, the internal energy density can be written as

$$u_{QGP} = \frac{3}{2\pi^2} \mu_q^4 - \mu_q^2 T^2 + \frac{37}{30} \pi^2 T^2 \quad (1.42)$$

and using 1.27 the QGP pressure is

$$p_{QGP} = \frac{1}{2\pi^2}\mu_q^4 - \frac{1}{3}\mu_q^2 T^2 + \frac{37}{90}\pi^2 T^4. \quad (1.43)$$

Applying the condition shown in equation 1.41, the critical temperature becomes

$$T_C = \left(\frac{90B}{37\pi^2} \right)^{1/4} = 145 \text{ MeV} \quad \text{for } \mu_q = 0, \quad (1.44)$$

which leads to the previous result where the chemical potential is zero. The additional information gained by using equation 1.40 is in case the chemical potential μ_q is non-zero. For $T \rightarrow 0$ it is possible to obtain the critical chemical potential for quarks in the QGP state as

$$\mu_q^C = (2\pi^2 B)^{1/4} = 0.43 \text{ GeV} \quad \text{for } T = 0. \quad (1.45)$$

This result can also be used to obtain the critical density of baryons for the QGP transition. Using the relation between the chemical potential and density of particles from [24] given by

$$n_k = \frac{N_k}{V} = \frac{\mu_k^3}{\pi^2} + \mu_k T^2. \quad (1.46)$$

Considering that $n_u + n_d = 2n_q = 3n_B$ (with $n_u = n_d$) and $\mu_u = \mu_d = \mu_q$, the critical density of baryons for $T = 0$ is

$$n_B = \frac{2}{3}n_q = \frac{2}{3\pi^2}\mu_q^3 = 0.72 \text{ fm}^{-3}. \quad (1.47)$$

The baryon density of nuclear matter at ordinary conditions is $n_B = 0.14 \text{ fm}^{-3}$ and the chemical potential is $\mu_B = 0.25 \text{ GeV}$, so the critical baryon density for QGP formation is about five times higher. At this baryonic density, the bag can no longer maintain quarks inside hadrons and the system changes to a state of free quarks.

Figure 1.7 presents the phase diagram obtained from equation 1.43 using the condition that the pressure should be equal to the bag pressure. Though this simple model provides basic aspects of the transition of the nuclear matter to the QGP, it is worth noting some further aspects of this phase transition, e.g. the baryonic chemical potential has a discontinuity along $T = 0$ until it reaches nuclear matter composed by protons and neutrons. With the increasing of the temperature this discontinuity moves towards smaller baryonic density [25].

Along the region $\mu_B < 0.25 \text{ GeV}$ (nuclear matter) and $T \rightarrow \infty$ there is a critical temperature where quarks confined in nuclear matter change to a deconfined state of quarks and gluons in a rapid crossover [26]. There are also reasons to infer that for a certain $T \neq 0$ and $\mu_B \neq 0$ there is a critical point where the phase transition is no longer a crossover, but a first order transition [27].

Figure 1.7: *Quark-Gluon Plasma phase diagram obtained by equation 1.43.*

Chapter 2

Signatures of the Quark-Gluon Plasma

“The deepest sin against the human mind is to believe things without evidence. Science is simply common sense at its best – that is, rigidly accurate in observation, and merciless to fallacy in logic.”

— Aldous Huxley

In 2000, the *Relativistic Heavy-Ion Collider* (RHIC) started operating at the *Brookhaven National Laboratory* (BNL). RHIC is a circular particle accelerator capable of performing proton-proton (pp) collisions up to 200 GeV in the center of mass and ion collisions (such as $Au+Au$) up to 200 GeV per nucleon.

One of the main goals of the experiments at RHIC is to find signatures of the QGP. In order to provide evidences for the existence of this state of matter a number of effects were measured such as flow (collective effects), particle energy loss and jet quenching.

Ten years later, in March 2010, the Large Hadron Collider (LHC) at CERN, in Geneva, started taking data of proton-proton collisions at $\sqrt{s} = 7$ TeV. In 2010 and 2011 the LHC also performed collisions using two beams of lead nuclei at $\sqrt{s_{NN}} = 2.76$ TeV. The purpose of taking heavy-ion data was the study of the Quark-Gluon Plasma, in which the ALICE experiment specializes.

Technical details of the LHC and the ALICE experiment will be covered ahead in chapter 4. This chapter will focus on three of the observables measured by both laboratories in heavy-ion collisions: ν_2 (elliptic flow), *Nuclear Modification Factor* (energy loss) and jets.

2.1 Geometry of Heavy-Ion Collisions

Nuclei collisions are characterized by their geometry. Figure 2.1 presents a collision where the projectiles are not head-on with an impact parameter $b > 0$. The lead nuclei in the left side of the picture are flat in order to illustrate the Lorentz contraction.

A centrality percentile from 0 to 100% can be defined by the impact parameter and it is commonly used in heavy-ion collisions experiments. In the extreme cases, $b = 0$ translates to a centrality of 0% and $b = r_1 + r_2$, where r_1 and r_2 are the radii of the two nuclei, to 100%. An expression for the centrality class can be written as

$$c = \frac{\int_0^b d\sigma/db' db'}{\int_0^\infty d\sigma/db' db'} = \frac{1}{\sigma_{AA}} \int_0^b \frac{d\sigma}{db'} db' \quad (2.1)$$

Figure 2.1: Schematic representation of a heavy-ion collision.

where σ is the total nuclear interaction cross section [28] and $d\sigma/db'db'$ is the impact parameter distribution.

This is a formal definition, in fact, the impact parameter is not experimentally accessible and the centrality is then defined as

$$c \simeq \frac{1}{\sigma_{AA}} \int_{N_{ch}^{THR}}^{\infty} \frac{d\sigma}{dN'_{ch}} dN'_{ch} \quad (2.2)$$

where $\frac{d\sigma}{dN'_{ch}}$ is the hadronic cross section for a given multiplicity (in some cases, charged particle multiplicity). In experiment, the cross section is replaced by the number of observed events. The integral in equation 2.2 is evaluated from a multiplicity threshold N_{ch}^{THR} . Below this threshold, the cross section of QED processes is large. These processes only affects the most peripheral collisions.

The Glauber model [28] connects the centrality percentile to the event multiplicity. The number of particles produced in a collision is related to the number of nucleons in the overlap area between the two nuclei in the collision. The right side of figure 2.1 illustrates that not all the nucleons take part in the collision. These nucleons are called spectators and those in the overlap region are the participants.

The centrality determination with ALICE detectors will be discussed in section 4.2.4.

2.2 Observables in Heavy-Ion Collisions

The main purpose of colliding heavy-ion nuclei is the study of nuclear matter in high temperatures and energy densities. As discussed in the previous chapter, in these conditions, hadrons should transit from a color confined state into a plasma of free quarks and gluons.

Figure 2.2 [29] presents a scheme of the space-time evolution of a heavy-ion collision starting in $z = t = 0$. The nuclei interact and deposit energy in the event region. The highly inelastic collision creates a large amount of (mostly light) quarks and gluons and the system enters in a pre-equilibrium phase.

Though not yet very well understood, the pre-equilibrium phase will give the initial conditions for the QGP phase when the system reaches local equilibrium. Parton energy loss and collective effects occur at this point and it is expected that matter behaves hydrodynamically.

Figure 2.2: Schematic representation of the QGP space-time evolution.

With the expansion of the system, the temperature drops and partons hadronize. These particles interact as a hadron gas with inelastic collisions until the chemical freeze-out is reached. At this moment, the dominant interactions are elastic scatterings. These scatterings stop at the kinetic freeze-out, after which there are only free hadrons.

2.3 Collective Effects

Considering the formation of the hot medium, the system thermalizes and particles present collective behavior that can be studied in heavy-ion experiments by measuring the distribution of particles in the $\eta - \phi$ phase space. Figure 2.3 [30] presents the geometry of a semi-central heavy-ion collision (left picture) with ellipsoidal geometry. The pressure in the x direction is higher than in y causing an azimuthal anisotropy in the momentum phase-space (right picture). Otherwise, if particles do not present collectivity, this momentum modulation will not be observed.

Figure 2.3: Schematics of a non-central heavy-ion collision (left) and anisotropic momentum distribution due to pressure gradients (right).

The azimuthal modulation is expressed quantitatively using a Fourier series

Figure 2.4: Measurement of ν_2 for charged particles (blue) and the average of D mesons species (black) measured by the ALICE Collaboration [31].

$$\frac{dN}{d(\phi - \Psi)} = \frac{1}{2\pi} \left\{ 1 + \sum_n 2\nu_n \cos(n[\phi - \Psi]) \right\}, \quad (2.3)$$

where Ψ is the orientation of the reaction plane (defined by the beam axis z and the impact parameter) with the x axis and ν_n is the n -th harmonic coefficient given by

$$\nu_n = \langle \cos(n[\phi - \Psi]) \rangle, \quad (2.4)$$

where the brackets denotes the average over all the selected particles. The second harmonic (ν_2) will quantify the magnitude of the elliptic contribution of the azimuthal modulation

Figure 2.4 presents the measurement of ν_2 as a function of the p_T of charged particles and the average of three species of D mesons (D^0 , D^+ and D^{*+}) measured by ALICE experiment at the LHC in Pb-Pb (30-50% centrality) collisions at $\sqrt{s_{NN}} = 2.76$ TeV [31]. For charged particles, dominated by charged pions, there is a clear evidence of elliptic flow.

The measurement of ν_2 of D mesons is an evidence of collective effects on heavy-flavour quarks and seems to follow the same behavior observed for light-flavoured particles.

2.4 Energy Loss

Energy loss due to particles traversing the plasma produced in relativistic heavy-ion collisions is one of the main phenomena studied in high-energy nuclear physics. Partons can lose energy by elastic collisions with other particles or gluon emission induced by the hot medium.

2.4.1 Collisional Energy Loss

A single parton can scatter multiple times in the QGP. It is then convenient to define a differential energy loss with respect to the space traveled by the parton inside the medium. Considering

a parton p traveling in the QGP, the differential energy loss should depend on the density of scattering centers (other quarks and gluons) in the QGP and the cross section of the parton to transfer energy to the medium in an elastic process. This can be described using [32]:

$$-\frac{dE_p}{dx} = \sum_{i=q,g} \int \rho_i(k) d^3k \int (1 - \cos \theta) \nu \frac{d\sigma_i^{elast}}{dq^2} dq^2 \quad (2.5)$$

where θ is the laboratory angle between the incident partons and $\nu = E_p^{in} - E_p^{out}$ with E_p^{in} , E_p^{out} being the energies of the incoming and outgoing partons and $\rho_i(k)$ is the density of particles of species i that can be a quark (q) or gluon (g) in the QGP with momentum k . An effective density of partons can be written as

$$\rho(k) = 2/3 \rho_q(k) + 3/2 \rho_g(k), \quad (2.6)$$

where the $2/3$ and $3/2$ factors take into account the nature of the quarks (fermions) and gluons (bosons).

The differential cross section for an energetic parton can be written as

$$\frac{d\sigma_i^{elast}}{dq^2} = \left(\frac{2}{3}\right)^{\pm 1} \frac{2\pi\alpha_s^2}{q^4} \quad (2.7)$$

where the sign depends on the parton type: $+$ for an incoming quark and $-$ for an incoming gluon.

Considering now $E_p^{in}, E_p^{out} \gg \nu$ and massless quarks, it is possible to use the Bjorken x relation for elastic processes

$$x = \frac{q^2}{2k \cdot \nu} = 1 \quad (2.8)$$

to obtain

$$q^2 = 2k \cdot \nu \simeq 2k\nu - k\nu \cos \theta \simeq 2k\nu(1 - \cos \theta) \quad (2.9)$$

and simplify equation 2.5 into

$$-\frac{dE}{dx} = \pi\alpha_s \left(\frac{2}{3}\right)^{\pm 1} \int \rho(k) \frac{d^3k}{k} \int \frac{dq^2}{q^2}. \quad (2.10)$$

In order to integrate in q^2 it is necessary to define a maximum and minimum value and there is more than one possible choice. The way done in [33] will be applied now and it is sufficient to illustrate some features of the collisional energy loss. The Debye mass is used as minimum momentum transfer $q_{min}^2 = m_D^2 \alpha_s T^2$ and for the maximum $q_{max} = \sqrt{4TE}$ with T the temperature of the QGP resulting in

$$-\frac{dE}{dx} = \pi\alpha_s \left(\frac{2}{3}\right)^{\pm 1} \int \rho(k) \frac{d^3k}{k} \ln \left(\frac{aE}{\alpha_s T} \right). \quad (2.11)$$

The previous choice of q_{min}^2 as the Debye mass also removed the k dependence from the logarithm in 2.11. The expression for the energy loss is then evaluated as

$$-\frac{dE}{dx} \approx 2\pi\alpha_s^2 T^2 \left(\frac{2}{3}\right)^{\pm 1} \left(1 + \frac{N_f}{6}\right) \ln\left(\frac{aE}{\alpha_s T}\right). \quad (2.12)$$

where $a = \mathcal{O}(1)$ is a constant and N_f is the number of flavours.

Equation 2.12 indicates that the collisional energy loss is proportional to T^2 , so it depends on the square root of the energy density in the medium. Also, there is a dependence of the strong interaction coupling α_s between the traversing parton and the QGP.

2.4.2 Radiative Energy Loss

Partons traveling through the Quark-Gluon Plasma can lose energy by gluon emission induced by the medium. This section will focus on quark energy loss and, at the end, the energy loss for the special case of heavy-flavour quarks.

A radiated gluon can rescatter multiple times inside the QGP. It is then of interest to define a mean free path of the emitted gluon λ and the length of the hot medium L . The energy loss of a quark with energy E can be evaluated from the differential gluon energy spectrum

$$-\Delta E = - \int \omega \frac{d^2 I}{d\omega dz}. \quad (2.13)$$

For the discussion, radiative energy loss will be treated as in [33] introducing the formation time of the radiation as

$$t_{form} \simeq \frac{\omega}{k_{\perp}^2} \quad (2.14)$$

where ω is the energy of the emitted gluon and $k_{\perp}$ is its transversal momentum with respect to the source quark. For $t_{form} \gg \lambda$ a group of scattering centers in the medium act coherently as a single one, so it will also be important to introduce the coherence length

$$l_{coh} \simeq \frac{\omega}{\langle k_{\perp}^2 \rangle_{coh}} \quad (2.15)$$

where the accumulated gluon transverse momentum assuming a random motion is given by

$$k_{\perp}^2 = N_{coh} \mu_D^2 \simeq \frac{\mu_D^2 L}{\lambda}. \quad (2.16)$$

with N_{coh} the number of the effective scattering centers in a medium with N sources such that $N_{coh} < N$.

From equations 2.15 and 2.16 one derives

$$l_{coh} \simeq \sqrt{\frac{\omega \lambda}{\mu_D^2}} \quad (2.17)$$

and N_{coh} can be written as

$$N_{coh} \simeq \sqrt{\frac{\omega}{\lambda \mu_D^2}} \equiv \sqrt{\frac{\omega}{E_{LPM}}}. \quad (2.18)$$

The parameter E_{LPM} is introduced here in analogy to the QED Landau-Pomeranchuk-Migdal (LPM) phenomenon. The number of effective scattering centers will also illustrate the different regimes in which the differential gluon energy spectrum is evaluated.

The differential gluon energy spectrum will be divided into three separated regimes.

- Bethe-Heitler regime ($\omega \leq \omega_{BH} \equiv E_{LPM}$)

For low energy gluons and large formation time $t_{form} \gg \lambda$, the scattering centers in the QGP act as independent sources of radiation ($N_{coh} \simeq N = L/\lambda$). The differential gluon energy spectrum per unit length is

$$\omega \frac{d^2 I}{d\omega dz} = \frac{\omega}{L} \frac{dI}{d\omega} \Big|_L \simeq \frac{\alpha_s N_{coh}}{\pi \lambda}. \quad (2.19)$$

- LPM regime ($\omega_{BH} < \omega < \omega_{fact} \simeq \mu_D^2 L^2 / \lambda$)

The scattering centers act as N_{coh} coherent sources ($1 < N_{coh} < N$) and $\lambda < l_{coh} < L$ in the LPM regime. The gluon spectrum is given by

$$\omega \frac{d^2 I}{d\omega dz} \Big|_{LPM} \simeq \frac{w}{l_{coh}} \frac{dI}{d\omega} \Big|_{l_{coh}} \simeq \frac{\alpha_s N_{coh}}{\pi l_{coh}} \simeq \frac{\alpha_s N_{coh}}{\pi} \sqrt{\frac{\mu_D^2}{\lambda \omega}}. \quad (2.20)$$

It is worth pointing out that by comparing equations 2.18, 2.19 and 2.20 a suppression of radiation given by the factor $\sqrt{E_{LPM}/\omega}$ emerges when the gluon energy rises from the Bethe-Heitler regime to the LPM.

- Factorization regime ($\omega_{fact} < \omega < E$)

In the factorization regime one finds the extreme case where only one effective scattering center is active ($N_{coh} \simeq 1$) and $l_{coh} > L$. The gluon spectrum is

$$\omega \frac{d^2 I}{d\omega dz} \Big|_{fact} \simeq \frac{\alpha_s N_{coh}}{\pi L}. \quad (2.21)$$

From the condition for ω in the factorization regime one notes that the three regimes are valid for a finite medium $L < L_{cr} = \lambda \sqrt{E/E_{LPM}}$.

The differential energy loss can be evaluated for the three regimes by integrating over ω with their proper condition for the integration limits.

$$- \frac{dE}{dz} \Big|_{BH} \simeq \frac{\alpha_s N_{coh} \mu_D^2}{\pi} \quad (Bethe - Heitler) \quad (2.22)$$

$$- \frac{dE}{dz} \Big|_{LPM} \simeq \frac{\alpha_s N_{coh}}{\pi} \sqrt{\frac{\mu_D^2 \omega_{fact}}{\lambda}} \simeq \frac{\alpha_s N_{coh} \mu_D^2}{\pi \lambda} L \quad (LPM) \quad (2.23)$$

$$- \frac{dE}{dz} \Big|_{fact} \simeq \frac{\alpha_s N_{coh} E}{\pi L} \quad (Factorization) \quad (2.24)$$

Interesting results appear when the relative energy loss $\Delta E/E$ is calculated by integrating in z

$$\frac{\Delta E}{E} \sim \frac{a \alpha_s L}{E} + \frac{b \alpha_s L^2}{E} + c \alpha_s, \quad (2.25)$$

where a , b and c are constants of $\mathcal{O}(1)$ and the terms in the right side of the expression are the contributions from the Bethe-Heitler, LPM and factorization regimes, respectively. This result shows that the loss of energy by gluon emission induced by the QGP grows with the square of the medium length L and the relative energy loss is smaller for high-energy quarks traveling in the QGP.

Now a special treatment will be given to heavy-flavour quarks. From [34] the distributions of soft gluons emitted by a heavy quark is given by

$$dP_{HF} = \frac{4\alpha_s}{3\pi} \frac{d\omega}{\omega} \frac{k_{\perp}^2 dk_{\perp}^2}{(k_{\perp}^2 + \omega^2 \theta_0^2)^2} \quad (2.26)$$

where θ_0 is defined as

$$\theta_0 \equiv \frac{M}{E}. \quad (2.27)$$

with M the mass and E the energy of the heavy quark.

The standard gluon distribution dP can be written in the approximation of small radiation angles $\theta \simeq k_{\perp}/\omega$ as

$$dP \simeq \frac{4\alpha_s}{3\pi} \frac{d\omega}{\omega} \frac{dk_{\perp}^2}{k_{\perp}^2} \simeq \frac{4\alpha_s}{3\pi} \frac{d\omega}{\omega} \frac{d\theta^2}{\theta^2}. \quad (2.28)$$

Comparing equations 2.26 and 2.28, the ratio is

$$\frac{dP_{HF}}{dP} = \left(1 + \frac{\theta_0^2}{\theta^2}\right)^{-2}. \quad (2.29)$$

The emission of gluon in small angles is suppressed by the factor obtained in 2.29 if the radiation angle θ is smaller than θ_0 and this is called *Dead-Cone Effect*. This factor also indicates that *bottom* quarks will lose less energy than *charm* quarks, since $M_{bottom} > M_{charm}$. In order to have a quantitative evaluation of the effect, a radiation amplitude $\overline{|R|^2}$ [35] is derived from 2.28 as

$$dP = \overline{|R|^2} \frac{d^3k}{2\omega(2\pi)^3}. \quad (2.30)$$

The radiation amplitude will be evaluated for the Feynman diagrams 2.5a (pre-emission) and 2.5b (post-emission), which are the *abelian* contribution. The last diagram 2.5a is the *non-abelian* contribution of three gluons vertex and will not be calculated. The *non-abelian* diagram was studied[35] and also indicates gluon emission suppression for heavy-quarks.

Figure 2.5: Feynman diagrams of one gluon radiation.

The pre-emission dominates in the angular region between $\theta_q \equiv |\vec{q}_\perp|/E$ and $\theta_c \equiv m^2/|\vec{q}_\perp|E = \theta_0^2/\theta_q$, for $\theta < \theta_c$ the post-emission dominates and in the region $\theta > \theta_q$ the two processes interfere.

Figure 2.6: Radiation amplitude as a function of θ for down (continuous line) and charm quark (dashed line). The parameters for the calculation are $p_i = (0, 0, 10)$ GeV, $p_f = (0.3, 0.2)$ GeV, $\mu_D = 0.5$ GeV for the quark and $\omega = 0.001$ GeV, $\phi = \pi/2$ for the radiated gluon. From [35].

Figure 2.6 presents the radiation amplitude $\overline{|R|}^2$ as a function of θ for the *abelian* contribution. The emission of soft gluons from the angular region $\theta > 10^{-2}$ rad is highly suppressed in comparison to light-quarks while the peak of emission for *down* quarks lies in the region $5 \cdot 10^{-4} < \theta < 7.5 \cdot 10^{-4}$ rad.

Heavy quarks traveling through the QGP should lose less energy than light quarks. A mass hierarchy of energy loss can be established as

$$\Delta E(\text{gluon}) > \Delta E(\text{light} - \text{flavour}) > \Delta E(\text{heavy} - \text{flavour}). \quad (2.31)$$

2.4.3 Nuclear Modification Factor

The most common observable used to quantify energy loss in high-energy collision experiments is the *Nuclear Modification Factor* (R_{AA}), which is expressed as

$$R_{AA} = \frac{1}{N_{\text{coll}}} \frac{d^2 N_{AA}/dp_T d\eta}{d^2 N_{pp}/dp_T d\eta} = \frac{1}{\langle T_{AA} \rangle} \frac{d^2 N_{AA}/dp_T d\eta}{d^2 \sigma_{pp}/dp_T d\eta} \quad (2.32)$$

where $d^2 N_{AA}/dp_T d\eta$ and $d^2 N_{pp}/dp_T d\eta$ are the particle (or jet) yield per unity of pseudorapidity and transverse momentum in nucleus-nucleus and proton-proton collisions, respectively, and N_{coll} is the number of binary collisions. The $\langle T_{AA} \rangle$ is the nuclear overlap function defined as N_{coll} over the inelastic pp collision cross section.

The R_{AA} is expected to be equal to 1 if there are no effects from the hot medium, i.e. no QGP formation. Energy loss due to the medium, for example, will be translated into R_{AA} smaller than unity.

Figure 2.7: Nuclear modification factor of D mesons (average of D^0 , D^+ and D^{*+}), charged pions and charged particles in 0-10% central Pb-Pb events at $\sqrt{s_{NN}} = 2.76$ TeV [36].

The nuclear modification factor for D mesons (average of D^0 , D^+ and D^{*+}), charge pions and charged particles [36] is presented in figure 2.7. Above 10 GeV/ c , the R_{AA} for D mesons starts rising (around 6 GeV/ c for light particles). Considering the energy loss models in this section it is expected the relative radiative energy loss to be smaller for high-momentum particles.

Despite the fact that the energy loss models are compatible with the rising trend in intermediate and high p_T , the parton mass dependence seems not to be confirmed in the presented data. The measurements of D -meson R_{AA} are still relatively new [36, 37] and the reason of its similarity with the results for light quarks is still a debate.

In order to fully understand the nuclear modification factor, one has to consider the different fragmentation processes that quarks undergo in the cases of light and heavy flavour particles.

2.5 Quark Fragmentation

High-energy collision experiments have shown that processes with a quark-antiquark pair production result in the observation of two (or more, as discussed in section 1.2) sprays of particles called jets.

There are two phenomena involved in this process of a quark ending up producing a jet. The first is the fragmentation of the initial quark into more quarks with less energy than the original. Then *hadronization* comes into scene and the resulting quarks will form colour singlets, i.e. hadrons.

The probability $\mathcal{D}$ for a hadron H with momentum p_H to be produced from a quark q with momentum p_q is given by the fragmentation function.

The *Lund* [38] fragmentation function is widely used in event generators [39]. The basis of this function is the *String* model [40] for fragmentation. A good illustration of the string model can be drawn from electron-positron annihilation going into a quark-antiquark pair. The quarks are produced back-to-back and the colour field connecting both particles can be described as a string with a effective potential

$$V(r) = kr, \quad (2.33)$$

where $k \approx 1 \text{ GeV/fm}$ is the string tension constant or linear energy density. The energy in the string gets higher with the moving away of the endpoint quarks. When the distance r between the quarks is of the order of $2m_T/k$, the string breaks up, creating a new quark-antiquark pair from vacuum excitation by the colour field.

This process can happen multiple times resulting in hadrons. The breaking up is controlled by a probability of quantum tunneling given by

$$f(z) \sim \exp\left(-\pi \frac{m_{q,\perp}^2}{k}\right) \quad (2.34)$$

where $m_{q,\perp}^2 = m_q^2 + p_{q,\perp}^2$ is the squared transverse mass of the newly produced quark.

The tunneling probability also introduces a transverse momentum in the new quark that has to be compensated by the new antiquark. One interesting feature of this probability is the mass dependence. It is easier to produce an *up* than a *strange* quark. Typically, if the probability of producing a light-quark u is normalized to 1, then the relative probability of producing $u : d : s : c$ quarks would be approximately $1 : 1 : 0.3 : 10^{-11}$. Heavy-flavour quarks are not expected to be produced from fragmentation processes.

Based on this discussion, the Lund fragmentation function is given by

$$f(z) = \frac{N}{z} z^{a_\alpha} \left(\frac{1-z}{z}\right)^{a_\beta} \exp\left(-\frac{bm_{H,\perp}^2}{z}\right) \quad (2.35)$$

where a_α , a_β and b are free parameters, N is a normalization parameter, $m_{H,\perp}$ is the hadron transverse mass and the light-cone momentum fraction z is

$$z = \frac{(E + p_{||})_H}{(E + p)_q}, \quad (2.36)$$

with $p_{||}$ the longitudinal momentum of the hadron with respect to the quark flight line. The parameters a_α and a_β are flavour dependent; the first depends on the flavour of the old quark and the last depends on the new one. However, usually these flavour dependent parameters are set equal and equation 2.35 is simplified as

$$f(z) = \frac{N}{z}(1-z)^a \exp\left(-\frac{bm_{H,\perp}^2}{z}\right). \quad (2.37)$$

The fragmentation function expressed in equation 2.37 is referred to as *symmetric Lund fragmentation function* and is commonly used to describe light-flavour fragmentation. Typical values of a and b are 0.3 and 0.8 respectively, which makes the distribution peaks around $z \approx 0.2$. A fragmentation that peaks in a region below $z \approx 0.5$ is commonly called *soft fragmentation* and this is a characteristic of light-flavoured hadrons.

The Lund fragmentation function for heavy-quarks has a correction implemented by Bowler[41] that takes into account the heavy endpoint quarks in the string model. However, for heavy-flavour fragmentation, the Peterson [42] fragmentation function will be discussed in this thesis.

The Peterson fragmentation function can be obtained from the kinematic argument that the amplitude of the process $Q \rightarrow H + q$ of a heavy-quark Q fragmenting into a hadron containing the heavy quark H and a light quark q is proportional to the inverse of $\Delta E = E_H + E_q - E_Q$. By using a mathematical trick for $m_Q^2/p^2 \ll 1$ and $m_H \approx m_Q$, ΔE can be simplified as

$$\Delta E = (m_Q^2 + z^2 p^2)^{1/2} + (m_q^2 + (1-z)^2 p^2)^{1/2} - (m_Q^2 + p^2)^{1/2} \propto 1 - \frac{1}{z} - \frac{\epsilon_Q}{(1-z)} \quad (2.38)$$

where p is the heavy-quark momentum, z is the energy fraction of the hadron with respect to the quark and ϵ is a parameter of the order of m_q^2/m_Q^2 , with m_q of the order of non-perturbative scale $\sim (1/2 \text{ to } 1)m_\rho \rightarrow m_q \sim (1/2 \text{ to } 1/8) \text{ GeV}$.

With this kinematic consideration and a factor $1/z$ for longitudinal phase space, the Peterson fragmentation function is given by

$$\mathcal{D}(z) = \frac{N}{z[1 - 1/z - \epsilon_Q/(1-z)]^2}. \quad (2.39)$$

The fragmentation function presented in equation 2.39 peaks around $z \approx 1 - 2\epsilon_Q$. For charm quarks, ϵ_c is typically 0.15, which results in a peak around 0.7 in z . Contrary to light-flavour, heavy-flavour has a *hard fragmentation* and this is corroborated by experiments[43, 44, 45, 46, 47, 48].

The measurement of fragmentation functions in high-energy experiments is performed with an important remark. Since the energy of quarks are not accessible, what is in fact measured is the energy (or momentum) of jets. In many cases the z (or $z_{||}$) variable is defined as

$$z_{||} = \frac{\vec{p}_{jet} \cdot \vec{p}_H}{|\vec{p}_{jet}|^2}, \quad (2.40)$$

which is the longitudinal momentum of the hadron with respect to the jet direction.

The momentum fraction $z_{||}$ as presented in equation 2.40 is the definition that will be used in this thesis.

2.6 Jets

Partons produced in the *hard scattering* of a high-energy collision commonly have a high virtuality Q^2 that decreases with the fragmentation (and radiation). After hadronization and, possibly, decay of some of the resulting hadrons, a group of collimated particles called *jet* is formed.

Experimentally, jets are defined as a group of particles inside a cone as presented in figure 2.8. The parameters involved in the definition of this cone depends on the algorithm used to reconstruct jets and will be discussed in section 5.3.

Figure 2.8: Jet pictured as particles inside a cone.

By summing the 4-momenta of the particles contained in the jet, one should reproduce information of the original parton such as energy, momentum and direction. As discussed in section 2.5, jets are also used to study quark fragmentation, which is of major importance in the understanding of the R_{AA} of hadrons, as mentioned in section 2.4.3.

Jets are also powerful tools to study parton collectivity and energy loss, bypassing the difficulties related to fragmentation and hadronization. In heavy-ion collisions, gluon radiation can be classified as *in-cone* radiation, when the angle of the emitted gluon is small enough to be inside the jet cone, or *out-of-cone*, when the angle is large and not reconstructed as part of the jet (also pictured in figure 2.8) and, consequently, reduces the value of the jet R_{AA} .

Figure 2.9 presents the measurement of jet nuclear modification factor done by ALICE experiment in two centrality classes [49]. For both centralities, $R_{AA} < 1$, indicating a strong suppression compatible with out-of-cone radiation. The predictions of JEWEL [50] and YaJEM [51] models are also displayed and are compatible with the measurement within uncertainties.

Figure 2.9: Jet nuclear modification factor in 0-10% (left) and 10-30% (right) centrality of Pb-Pb collisions at $\sqrt{s} = 2.76$ TeV with respect to the same measurement done in pp collisions. Data is compared to JEWEL (black line) and YaJEM (green dashed line) models [49].

Chapter 3

Heavy Flavour

“To the makers of music – all worlds, all times.”

— Voyager Golden Record inscription

Charm quarks, will be the entities used to probe the Quark-Gluon Plasma in this thesis. The experimental and theoretical advantages of using heavy quarks will be discussed in this chapter. Furthermore the production mechanisms and previous measurements will also be presented.

3.1 Production Mechanisms

Heavy-flavour quarks are characterized by their large mass in comparison to the QCD scale $\Lambda \approx 200$ MeV. Hence, it is commonly stated that $m_q \ll \Lambda \ll m_Q$, where m_q is the mass of the light-flavour quarks u , d and s and m_Q is the mass of c , b and t heavy quarks.

The momentum transfer involved in heavy quark production has to be at least equal to two times the quark mass m_Q . At these energy scales ($Q^2 \geq m_Q^2 \gg \Lambda^2$) the strong coupling constant $\alpha_s < 1$ and the QCD processes can be treated perturbatively in powers of α_s .

Due to the large mass, the production of heavy quarks is expected to be involved in the hard scattering. In hadronic collisions, heavy-flavour is produced in the early stages. The *Leading-Order* (LO) contribution in perturbative QCD, proportional to α_s^2 , will lead to Feynman diagrams of $q\bar{q} \rightarrow Q\bar{Q}$ (*annihilation*) and $gg \rightarrow Q\bar{Q}$ (*gluon fusion*). The *Next-to-Leading-Order* (NLO) corrections are diagrams proportional to α_s^3 and are represented by $q\bar{q}g \rightarrow Q\bar{Q}$, $qg \rightarrow Q\bar{Q}g$ and $gg \rightarrow Q\bar{Q}g$.

The production mechanisms can be classified according to the number of heavy quarks in the final state of the hard subprocess:

- *Flavour Creation*

This is the LO contribution and, as mentioned above, can be represented by Feynman diagrams as annihilation (figure 3.1a) or gluon fusion (figure 3.1b). The quark-antiquark pair is created back-to-back in order to conserve the momentum, though radiated gluons can change the kinematics (figure 3.1c). Flavour creation is also commonly called *pair creation*. Since all production mechanisms produce quarks in pairs, in this thesis the term *flavour creation* will be adopted.

Figure 3.1: Flavour Creation diagrams. Two heavy quarks in the hard subprocess.

- *Flavour Excitation*

When a parton scatters against an off-mass-shell heavy-quark of the parton distribution and put it on mass shell, it is called flavour excitation. The process $qg \rightarrow Q\bar{Q}q$ presented in figure 3.2b is an example of this mechanism. Flavour excitation can have initial and final state partons similar to flavour creation, e.g. $gg \rightarrow Q\bar{Q}g$ (figure 3.2a). However, in flavour creation, the quark-antiquark pair is produced in the more virtual subprocess (hard scattering), not in the shower involved in the hard interaction.

Figure 3.2: Flavour Excitation diagrams. One heavy quark in the hard subprocess.

- *Gluon Splitting*

In this process, a gluon breaks up into a quark-antiquark pair. The shower can occur in the initial or final state. Heavy-quark production in the final state (figure 3.3b) is the dominant process because of the maximum virtuality restriction of gluons in the initial state. In initial state production the pair can also come from the parton distribution being put on mass shell by the interaction with the gluon and later emit a gluon that will take part in the hard scattering (figure 3.3a). This is, essentially, a flavour excitation process, but since no heavy quark is in the hard subprocess, it is classified as gluon splitting.

Figure 3.3: Gluon Splitting diagrams. No heavy quarks in the hard subprocess.

Heavy-flavour production in hadron collisions is strongly dependent of the *Parton Distribution Function* (PDF), which describes the distribution of partons inside hadrons for a given energy scale. The evolution of the PDFs is such that for sufficiently high energy scales, not only gluon and the valence quarks (u and d) in hadrons, e.g. protons, are relevant, *strange* and, for increasing energies, *charm*, *bottom* and *top* also have to be taken into account.

The production cross section of a heavy hadron in hadron collisions is then dependent of perturbative (heavy-quark production) and non-perturbative (fragmentation and hadronization) processes. J. C. Collins demonstrated [52] that the calculation of the heavy-hadron cross section can be factorized due to the different time scales involved in each process takes place. The differential cross section for a heavy hadron production can be written as

$$\frac{d\sigma^{pp \rightarrow HQ+X}}{d^3\vec{p}_H} = \int \frac{d^3\vec{p}_Q}{E_Q} E_Q \frac{d\sigma^{pp \rightarrow Q\bar{Q}+X}}{d^3\vec{p}_Q} \int dz \mathcal{D}_Q^H(z) \delta^{(3)}(\vec{p}_H - z\vec{p}_Q) \quad (3.1)$$

where $\mathcal{D}_Q^H(z)$ is the fragmentation function of a quark Q to originate a hadron H that carries a fraction z of the quark energy. The production cross section of a heavy quark-antiquark pair $E_Q \frac{d\sigma^{pp \rightarrow Q\bar{Q}+X}}{d^3\vec{p}_Q}$ can be factorized as

$$E_Q \frac{d\sigma^{pp \rightarrow Q\bar{Q}+X}}{d^3\vec{p}_Q} = \sum_{i,j} \int dx_1 dx_2 f_i^p(x_1, \mu_F^2) f_j^p(x_2, \mu_F^2) E_Q \frac{d\hat{\sigma}^{ij \rightarrow Q\bar{Q}}}{d^3\vec{p}_Q} \quad (3.2)$$

where the sum is carried out over all partons i, j with fraction x_1 and x_2 of the momentum of the hadron in the beam. The PDF f_i^p is calculated for the factorization scale μ_F . The partonic cross section for the production of a heavy quark-antiquark pair is represented by $\hat{\sigma}^{ij \rightarrow Q\bar{Q}}$ and is called Wilson coefficient. These coefficients are the hard-scattering amplitudes from the processes calculated by perturbative QCD that were discussed in the beginning of this section.

Relying on the factorization and the production mechanisms obtained using perturbative QCD, the contribution of each process to the charm production cross-section for charm (figure 3.4a) and bottom (figure 3.4b) in pp collisions are presented[53]. At the LHC energies, flavour excitation is the dominant process for charm and bottom production. For charm, gluon splitting is considerably more relevant than flavour creation, which is not true for beauty. The total cross section presented in the figures 3.4a and 3.4b represents the sum over the three processes, so no non-perturbative processes are taken into account.

Figure 3.4: Charm (a) and bottom (b) production cross sections in pp collisions.

3.2 Heavy-Flavour in Heavy-Ion Collisions

Heavy quarks are very useful tools to study cold and hot nuclear effects. Due to their large masses, heavy-flavour particles are produced in the early stages of hadronic collisions and will be affected by nuclear matter.

Considering the *dead-cone effect* discussed in section 2.4.2, it is expected that heavy quarks lose less energy than light quarks or gluons. Furthermore, it is also expected that bottom will lose even less energy than charm since the former has a larger mass. This mass dependent energy loss is an important property of the medium that can be tested with open heavy flavour.

Until now, the fragmentation function for heavy quarks was never measured in heavy-ion collisions. It is still unknown if initial or final state nuclear matter effects can modify the fragmentation. Factorization states that fragmentation is independent of the PDFs, so no modification in the fragmentation functions is expected.

For a better understanding of the modifications on observables due to the Quark-Gluon Plasma (final state effects), it is of major importance to evaluate cold nuclear matter (initial state) effects that originate from the fact that a nucleus with N nucleons will not behave as N unbound nucleons, changing the nuclear PDF.

Cold nuclear matter can be quantified by the ratio between the nuclear structure function, for instance F_2 , of the nucleus over that for the nucleon[54] as

$$R_{F_2}^A(x, Q^2) = \frac{F_2^A(x, Q^2)}{A F_2^{\text{nucleon}}(x, Q^2)}. \quad (3.3)$$

The ratio presented in equation 3.3 defines four different regions of effects:

- $x \gtrsim 0.8$: *Fermi Motion* $\rightarrow R_{F_2}^A > 1$

In this region the measured ratio is higher than unity due to the fact that the nucleons are not static inside the nucleus. This motion changes the experimental Bjorken- x of the partons

that is usually defined in the approximation of a static nucleus.

- $0.25 \lesssim x \lesssim 0.8$: EMC $\rightarrow R_{F_2}^A < 1$

Though the EMC effect is clearly observed in experiment, its theoretical interpretation is somewhat ambiguous. The measurement comparison between iron nucleus and deuterium showed that the Fermi motion is smaller for heavier nuclei [55].

- $0.1 \lesssim x \lesssim 0.25$: Antishadowing $\rightarrow R_{F_2}^A > 1$

Antishadowing can be understood as the fusion of two partons of small x interacting with each other, in which case the x probed is higher.

- $x \lesssim 0.1$: Shadowing $\rightarrow R_{F_2}^A < 1$

The ratio $R_{F_2}^A$ below unity can be understood as the nucleons in the surface of the nucleus casting a shadow over the nucleons inside it. The shadowing effect is larger for heavier nuclei.

Observables such as the nuclear modification factor can be sensitive to both kinds of nuclear effects. In order to disentangle the contribution from each source, measurements of hadron collisions in which the formation of the QGP is not expected, such as in p-Pb collisions, are of great interest.

Collisions of protons against nuclei can be used to evaluate cold nuclear effects. In the ALICE experiment at the LHC, observables are measured in proton-proton (usual baseline) and proton-lead collisions in order to make comparisons. For instance, the nuclear modification factor R_{pPb} of particles in function of the particle transverse momentum.

Figure 3.5: Average R_{pPb} of D mesons in pPb collisions at $\sqrt{s_{NN}} = 5.02$ TeV compared with models. Statistical (bars) and systematic (boxes) uncertainties are also presented [56].

The nuclear modification factor measured in proton-nucleus collisions should increase or decrease in the same direction of $R_{F_2}^A$, so the shadowing effect is translated to $R_{pA} < 1$. The measurement of the average R_{pPb} for D mesons (D^0 , D^+ and D^{*+}) is presented in figure 3.5 [56] and is compatible

with models. However, due to the relatively large uncertainties, one can not state that cold nuclear effects were observed.

Chapter 4

ALICE Experiment

“It doesn’t matter how beautiful your theory is, it doesn’t matter how smart you are. If it disagrees with experiment, it’s wrong.”
— Richard Feynman

ALICE stands for *A Large Ion-Collision Experiment* and it is one of the four main experiments at the *Large Hadron Collider*. In the beginning of this chapter the LHC and ALICE experiment will be presented, then some experimental aspects such as ALICE subsystems, centrality determination, vertices reconstruction, tracking and particle identification will be discussed.

4.1 Large Hadron Collider

The Large Hadron Collider is a circular collider with 27 km of circumference built in France and Switzerland at CERN, the European Organization for Nuclear Research.

Two beams of particles are accelerated in opposite directions to close-to-light speed inside the two LHC rings. The particles collide at specific points (interaction points) where the experiments are built, as presented in figure 4.1 [57].

The proton beam starts at Linac 2, where electrons are removed from hydrogen atoms that are accelerated; the hydrogen atoms are accelerated at the energy of 50 MeV. From the linear accelerator, the proton beam goes through three circular pre-accelerators. The *Booster* takes the beam to 1.4 GeV and sends it to the *Proton Synchrotron* (PS). There the beam gets to the energy of 25 GeV. The last accelerator before the LHC is the *Super Proton Synchrotron* (SPS), used in the past to produce proton-antiproton collisions. It pushes the energy to 450 GeV and injects the beam inside the Large Hadron Collider.

Since 2010, when it became fully operational, the LHC is able to successfully produce proton-proton collisions at $\sqrt{s} = 2.76, 5.02, 7$ and 13 TeV. The machine is also capable of producing heavy-ion collisions, which is of major importance in the study of the Quark-Gluon Plasma. Pb-Pb collisions were performed in 2010, 2011 ($\sqrt{s_{NN}} = 2.76$) and 2015 ($\sqrt{s_{NN}} = 5.02$). Proton-lead (pPb) collisions at $\sqrt{s_{NN}} = 5.02$ (2013 and 2016) and 8 TeV (2016) were also produced in order to investigate initial state effects.

Figure 4.1: Schematic representation of the Large Hadron Collider complex.

4.2 ALICE

The ALICE experiment is a general-purpose detector optimized for heavy-ion collisions. The main subject explored by ALICE is the study of the Quark-Gluon Plasma.

The ALICE Collaboration is composed by over 1900 scientists from over 30 countries. The detector is located 60 meters underground at interaction point 2 in Saint-Genis Pouilly, France. The overall dimension of the detector is $16 \times 16 \times 26 \text{ m}^3$ with a weight of ~ 10000 tons.

The ALICE detector [58] (figure 4.2) is composed of a central barrel used to measure hadrons, electrons and photons and a forward system dedicated to the study of muons. There are other subsystems used for event triggering and characterization, e.g. multiplicity is determined by V0 detector. The central part is covered by the L3 magnet that is usually set to produce a magnetic field of 0.5 Tesla.

The central barrel is composed of the *Inner-Tracking System* (ITS), the detector closest to the interaction point and, together with the *Time-Projection Chamber* (TPC), is responsible for vertex determination and the reconstruction of charged particle tracks. The TPC and *Time-of-Flight*

Figure 4.2: Picture of the ALICE detector and subsystems.

(TOF) provide particle identification (PID) determined by energy loss (TPC) or the different times that a particle, for a given momentum, takes to travel from the moment of the collision until it hits the detector.

The Transition Radiation Detector (TRD), used for the study of electrons, the Electro-Magnetic Calorimeter (EMCal-DCal), which provides electron and photon detection, the *High-Momentum PID* (HMPID), that performs particle identification using Cherenkov radiation, and the Photon Spectrometer (PHOS), dedicated to photon detection, are also part of the central barrel.

The following subsections will be dedicated to the most relevant detectors for this thesis. Important features of the detector in the context of this analysis, such as particle identification, tracking, centrality determination, etc, will be highlighted. For a better description of the other subsystems, reference [58] can be used.

4.2.1 Inner-Tracking System

The Inner-Tracking System is part of the central barrel and it is composed by six layers with full azimuthal coverage in three different configurations as presented in figure 4.3.

Figure 4.3: The six layers of the Inner-Tracking System.

- *Silicon Pixel Detector (SPD)*

The two innermost layers of the ITS. The two main tasks of this silicon detector are the determination of the primary vertex and impact parameter of secondary particles coming from weak decays of strange, charm or beauty hadrons. In heavy-ion collisions, it is expected that the number of tracks passing through the SPD is around 80 tracks per cm^2 , hence the necessity of high granularity of these layers that provide two-track resolution of $100 \mu\text{m}$ in $r\phi$.

- *Silicon Drift Detector (SDD)*

Two intermediate layers that provide two particle energy loss samples for particle identification. The expected density is 7 tracks per cm^2 .

- *Silicon Strip Detector (SSD)*

The outermost layers have the task of connecting the reconstruction of tracks from the ITS to the TPC. They also provide two more samples of dE/dx for the ITS particle identification.

Layer	Type	r (cm)	$\pm z$ (cm)	η	Spatial precision (μm)	
					$r\phi$	z
1	Pixel	3.9	14.1	1.98	12	100
2		7.6	14.1	1.4		
3	Drift	15.0	22.2	0.9	38	28
4		23.9	29.7	0.9		
5	Strip	37.8	43.1	0.9	20	830
6		42.8	48.9	0.9		

Table 4.1: Dimensions and two-track resolution of the ITS layers.

The dimensions and spatial precision information is summarised in table 4.1. The hadronic decay length of D mesons is of the order of $\sim 100 \mu\text{m}$. Therefore, these mesons are expected to decay inside the ITS and the detector precision will be exploited for the secondary vertex reconstruction.

4.2.2 Time-Projection Chamber

The Time-Projection Chamber is the main tracking detector and responsible for the reconstruction of charged particles. A schematic picture of the TPC is presented in figure 4.4. It consists of a 510 cm long cylinder with inner radius of 85 cm and outer radius of 250 cm. The detector is divided in half in the center by a high voltage electrode that generates an uniform electric field in the longitudinal direction pointing from the endplates to the center. It has full azimuthal coverage and $|\eta| < 0.9$ in pseudorapidity for tracks that do not go out of the detector through the endplates and are therefore fully contained in the TPC.

Figure 4.4: The six layers of the Inner-Tracking System.

The detector is filled with a combination of Ar (88%) and CO₂ (12%) gases that ionize when a charged particle passes through, leaving a trace of ions and electrons. Due to the electric field, the electrons move longitudinally to the endplates where they are collected. The bending of the trace depends on the particle momentum. Moreover, the density of electrons reaching the endplates depends on the particle momentum and species, which allows not only momentum determination but also particle identification.

The average particle energy loss per unity of path length in the TPC can be described by the Bethe-Bloch equation

$$\left\langle \frac{dE}{dx} \right\rangle = \frac{4\pi N e^4 Z^2}{m_e \beta^2} \left[\ln \left(\frac{2m_e \beta^2 \gamma^2}{I} \right) - \beta^2 - \frac{\delta(\beta)}{2} \right] \quad (4.1)$$

where $m_e = 0.51$ MeV is the electron rest energy, Z is the charge of the particle traversing the gas, N is the density of electrons in the trace, e is the elementary charge, β is the velocity of the particle and I is the mean excitation energy, which is calculable for simple gases, but, in general, it is a fit parameter. The term $\delta(\beta)$ is a correction to take into account the atom polarization in the region the particle is passing through and depends on the gas.

Particle identification with the TPC can be performed on a track-by-track basis in the region where the expected energy losses do not overlap. The energy loss in arbitrary units is presented in

figure 4.5 [59].

Figure 4.5: Average energy loss per unit of path length (arbitrary units) as a function of the transverse momentum over the particle charge.

Kaons can be identified without overlap with expected energy losses of other particles up to 500 MeV/c and pions up to 600 MeV/c in momentum. For particle identification at higher momentum the TPC energy loss information is combined with the information from TOF.

The track reconstruction in the TPC is based on the Kalman filter method [60]. The electrons collected in the endplates are converted into space-point positions. The reconstruction is done in two steps, first the tracks are reconstructed starting from the outermost space points and going inward and, in the second stage with the information from the previous reconstruction, it is done outward.

Inward reconstruction uses different types of seeds for different hypotheses. Considering the particle originated in the primary vertex, the vertex position is used as a constraint. The other seed is not constrained to the vertex, with the hypothesis that it is a secondary particle from a decay or produced in the detector material. The reconstruction ends with the propagation of the track to the ITS points.

With the track parameter estimates, the outward reconstruction is done removing the space points with high values of χ^2 . The Kalman filter ends with the propagation of the tracks to the TRD and TOF points. The TPC is capable of reconstructing charged particles from 0.1 GeV/c up to 100 GeV/c.

4.2.3 Time-of-Flight

The Time-of-Flight (TOF) detector consists of 1593 Multigap Resistive Plate Chambers (MRPC) divided in 18 supermodules arranged to give full coverage in azimuth. Each supermodule contains

5 modules in line with a total of 9 m long giving a coverage of $|\eta| < 0.9$ in pseudorapidity.

Particle identification with TOF is performed based on the variable

$$\delta_i = \frac{t_{TOF} - t_{ev} - t_{exp}(m_i, p, L)}{\sigma_{tot}(m_i, p, t_{ev})} \quad (4.2)$$

where t_{TOF} is the time marked when a particle $i = \pi, K, e, etc$ hits the detector, t_{ev} is the time of the collision given by the T0 detector and t_{exp} is the expected time of flight of a particle i for a given momentum p and path length L . The total uncertainty σ_{tot} is the result of a quadratic sum of TOF (~ 80 ps), t_{ev} and tracking reconstruction uncertainties.

The variable δ can be used to implement $n\sigma$ cuts, which is the deviation expressed in multiples of the uncertainty from the expected time of flight of the particle. The performance is presented in figure 4.6 [59] (β as a function of particle momentum), where good separation (3σ) between Kaon and Pion can be seen up to 2.5 GeV/c.

Figure 4.6: *TOF β in function of particle momentum.*

4.2.4 VZERO

The VZERO detector (or simply V0) consists of two scintillators located on each side of the interaction point. The V0A and V0C are positioned 340 cm and -90 cm from the vertex respectively. Each scintillator has 32 counters arranged in 4 rings, as presented in figure 4.7 [61].

A minimum bias trigger is provided by a combination of the V0 and ITS detectors. In order to separate heavy-ion collisions from events in which the beam collides with the gas in the LHC pipe, ALICE works with a synchronization of hits in the V0A and V0C and the LHC clock (25 ns). If the time of the bunch crossing is considered $t = 0$, the hits in the V0C should occur at $t \sim 3$ ns and in V0A at $t \sim 11$ ns. These time windows are different for beam-gas interaction coming from the V0A or V0C sides. The minimum bias characterization with the V0 is also combined with the SPD layer of the ITS, since they cover different acceptance regions.

Figure 4.7: V0A and V0C schematic rings.

Collision centrality is also estimated with the V0 detector. The energy deposited by particles in the V0 is used to estimate the number of primary charged particles in a given pseudorapidity range.

Figure 4.8: Number of events as a function of the V0 amplitude. The shaded areas mark regions attributed to centrality percentiles. The fit (in red) is performed using Glauber model.

As discussed in section 2.1, the connection between the multiplicity, or V0 amplitude, and the centrality is done by Glauber model. The centrality determination is presented in figure 4.8[61] and demonstrates that the Glauber fit successfully describes V0 amplitude.

Chapter 5

Analysis

“To achieve great things, two things are needed: a plan and not quite enough time.”

— Leonard Bernstein

The analysis was divided into two phases. The first phase the methods were developed and tested in order to establish a consistent framework. Also, in the first phase the focus was the measurement of the transverse momentum spectrum of jets produced by charm quarks.

In order to identify jets coming from a charm quark, a technique where jets are reconstructed with the D-meson candidate instead of the daughter tracks was applied. Later, the D-meson signal is extracted using two different methods, which leads to the yields of jets containing D mesons.

The D-meson contribution from bottom decays (feed-down) was subtracted and only D mesons coming from charm quark were considered (prompt D mesons). The open charm species used in this work was the D^0 .

This thesis describes the first phase of the analysis. The second phase consists in the measurement of the jet momentum fraction ($z_{||}$) and will not be discussed in this thesis.

Corrections and systematics for phase one will also be discussed in this chapter.

5.1 Event Selection

This work analysed the data collected in the late 2015 with ALICE at the LHC. These data contain information on collisions of lead nuclei (Pb-Pb collisions) at $\sqrt{s_{NN}} = 5.02$ TeV and the interest is on those most central events, i.e. events between 0% and 20% centrality.

In 2015 no centrality trigger was used and the data is minimum bias. The centrality distribution is presented in figure 5.1. The analysis also requires that the primary vertex is found in a longitudinal region $|z_{vtx}| < 10$ cm.

During 2015, 2.7×10^8 Pb-Pb events were collected and 1.3×10^7 (0-20% most central) were selected for analysis in this work. This dataset is labeled **LHC15o**.

Figure 5.1: Number of events as a function of the event centrality determined by V0 detector.

5.2 D-Meson Reconstruction

ALICE is capable of fully reconstructing D^0 , D^+ , D^{*+} and D_s^+ through hadronic channels. This work will focus on D^0 and this choice is made based on technical and statistical reasons. D^0 is the lightest charmed meson hence it is the most abundant meson carrying a charm quark produced at the LHC. Other charmed mesons, e.g. D^{*+} , have additional difficulties to be reconstructed in a heavy-ion environment due to a soft pion that is part of a three body decay and, in a Pb-Pb collision with a large number of pions, this results in a large background that has to be treated carefully.

The reconstruction of the D^0 includes the antiparticle $\bar{D}^0$, though it will always be referred only as D^0 unless explicitly remarked. The lightest charmed meson has a hadronic decay channel given by

$$D^0 \rightarrow K^- + \pi^+ \quad BR = 3.89\%. \quad (5.1)$$

The Branching Ratio (BR) indicates the percentile of a particle decaying in a given channel and has to be taken into account in order to obtain the full D^0 production.

The reconstruction strategy consists of combining two tracks of opposite charges that were identified as kaon or pion. These tracks also have to pass the following criteria:

- reconstructed with a minimum of 70 clusters in the TPC;
- reconstructed with a reduced $\chi^2 < 4$;
- 2 hits in the ITS, whereas at least one of them has to be in the SPD;
- $|\eta| < 0.8$.

The secondary vertex position is then determined with these two particles and the D^0 kinematics (mainly p_T , η and ϕ) is calculated. These reconstructed D mesons are called *candidates*.

The decay shown in 5.1 occurs via weak interaction and has a proper decay length of $c\tau \sim 123 \mu m$. This displacement of a few hundred micrometers from the primary vertex is exploited by ALICE capability of reconstructing tracks coming from a secondary vertex.

Based on the topology of the decay and, in order to suppress combinatorial background, the candidates have to pass a set of cuts that is summarized in table 5.1.

$p_T^{D^0}$ (GeV/c)	0-0.5	0.5-1	1-2	2-3	3-4	4-5	5-6	6-8	8-12	12-16	16- ∞
ΔM_{D^0}	0.4	0.4	0.4	0.4	0.4	0.4	0.4	0.4	0.4	0.4	0.4
DCA (cm)	0.04	0.04	0.04	0.025	0.025	0.025	0.025	0.03	0.035	0.04	0.04
$\cos \theta^*$	0.8	0.8	0.8	0.8	0.8	0.8	0.8	0.8	1	1	1
p_T^K (GeV/c)	0.4	0.4	0.4	0.7	0.7	0.7	0.7	0.7	0.7	0.7	0.7
p_T^π (GeV/c)	0.4	0.4	0.4	0.7	0.7	0.7	0.7	0.7	0.7	0.7	0.7
d_0^K (cm)	0.1	0.1	0.1	0.1	0.1	0.1	0.1	0.1	0.1	0.1	0.1
d_0^π (cm)	0.1	0.1	0.1	0.1	0.1	0.1	0.1	0.1	0.1	0.1	0.1
$d_0^K \cdot d_0^\pi$ (10^{-4} cm ²)	-4.3	-4.3	-4.3	-4.5	-3.6	-2.7	-2.1	-1.4	-0.5	-0.1	-0.1
$\cos \theta_{point}$	0.85	0.85	0.85	0.95	0.95	0.95	0.92	0.88	0.85	0.83	0.82
$\cos \theta_{point}^{XY}$	0.995	0.995	0.995	0.998	0.998	0.998	0.998	0.998	0.998	0.998	0.998
$L_{XY}/\sigma_{L_{XY}}$	8	8	8	7	5	5	5	5	5	8	6

Table 5.1: Kinematical and topological cuts for D^0 candidates in Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV.

The D^0 decay is presented schematically in figure 5.2. The pointing angle θ_{point} is the angle formed by the momentum direction of the reconstructed D meson and its flight-line defined by the primary and secondary vertices. This angle is expected to be small for D meson created in the primary vertex. The same angle projected in the x-y plane (θ_{point}^{XY}) is also applied in the selection.

The impact parameter d_0 is the distance of closest approach between the primary vertex and the backward projection of the reconstructed daughter tracks of the D^0 decay. As presented in figure 5.2, the daughters should have opposite impact parameter signs, so the product of the impact parameters of the daughter tracks is also used as selection criterion.

The DCA is the distance of closest approach between the two daughter tracks. The θ^* angle is formed between the kaon in the D^0 rest frame and the boost direction. L_{XY} is the decay length in the x-y plane. The difference between the invariant mass calculated with the daughter tracks and the D^0 mass (ΔM_{D^0}) is applied to reject candidates too far from the right mass region $m_{D^0} = 1864.84 \pm 0.05$ MeV [18].

Figure 5.2: Schematics of the D^0 decay by the hadronic channel.

The cut values in table 5.1 were chosen aiming the maximum significance $\mathcal{S}$ of the signal given by

$$\mathcal{S} = \frac{S}{\sqrt{S+B}}, \quad (5.2)$$

where S is the signal and B is the background taken from the functions used to fit the invariant mass distribution of the candidates

$$M_{D_0} = \sqrt{(E_K + E_\pi)^2 - (\vec{p}_K + \vec{p}_\pi)^2}, \quad (5.3)$$

with E and $\vec{p}$ the energy and momentum of the daughter particles.

Usually the signal extraction is performed in intervals of D-meson transversal momentum, however in this work the intervals of jet momentum (p_T^{jet}) or momentum fraction ($z_{||}$) has to be considered. This will be discussed later in section 5.5. The following sections will present the jet reconstruction strategy adopted to select heavy-flavour jets.

5.3 Jet Reconstruction

Jets in this work were reconstructed by FastJet [62], which is a C++ package containing jet finders and other analysis tools. For jets the *anti* - k_t algorithm was applied. This algorithm is *infrared-safe*, which means that it is robust against soft radiation, and *collinear-safe*, i.e. not sensitive to particle splitting.

The *anti* - k_t algorithm uses the momentum of the particles to clusterize them in jets. The algorithm works in the following way:

- calculate all variables d_{ij} and $d_{i\text{Jet}}$ with the i and j particles in the event, where

$$d_{ij} = \text{minimum} \left(\frac{1}{p_{T,i}}, \frac{1}{p_{T,j}} \right) \frac{\Delta R_{ij}^2}{R^2} \quad (5.4)$$

and

$$d_{i\text{Jet}} = \frac{1}{p_{T,i}}, \quad (5.5)$$

where

$$\Delta R_{ij}^2 = (\phi_i - \phi_j)^2 + (\eta_i - \eta_j)^2 \quad (5.6)$$

and R is the jet *resolution parameter*;

- find the smallest value between all the d_{ij} and $d_{i\text{Jet}}$ variables. If d_{ij} is smaller than $d_{i\text{Jet}}$, particles i and j are removed from the list and merged in a new particle k (in fact, a jet), otherwise, $d_{i\text{Jet}}$ is declared as a jet and is removed from the list of remaining particles;
- go back to the first step and calculate the variables with the remaining particles;
- stop when there are no more particles left in the list.

The algorithm basically starts the jet clusterization by the hard seeds, i.e. particles with high transverse momentum. The shape of the jets in the $\eta - \phi$ phase space is close to a circle. The resolution parameter R regulates the size of jets in the phase space, the probability of a jet to be formed with a radius larger than R decreases drastically.

The algorithm also requires a systematic way to merge the properties of two particles. This analysis uses what is called *p_T scheme*, where the merging of particles follow the recipe

$$p_{T,k} = p_{T,i} + p_{T,j}, \quad (5.7)$$

$$\phi_k = (p_{T,i}\phi_i + p_{T,j}\phi_j)/p_{T,k}, \quad (5.8)$$

$$\eta_k = (p_{T,i}\eta_i + p_{T,j}\eta_j)/p_{T,k}. \quad (5.9)$$

This analysis used only charged jets, i.e. jets clusterized only with charged tracks. In order to guarantee that the full jet cone is contained in the TPC acceptance, a requirement of $|\eta^{jet}| < 0.9 - R$ in pseudorapidity for charged jets was applied.

5.3.1 Background Subtraction

Jet reconstruction in heavy-ion collision is challenging due to the large number of particles produced in such collisions. A large amount of these particles do not come from a hard scattering, but from soft processes, such as particles created in the QGP.

These particles create a background in the environment in which jets coming from hard interaction are being reconstructed. Consequently, correlated particles from a parton fragmentation, i.e. a signal jet, will be reconstructed together with uncorrelated particles.

Another possibility is that a jet can be clusterized only with uncorrelated particles, which is called *combinatorial jet*, or combinatorial background. The last case that will be mentioned here

is when a signal jet is split into two (or more) jets due to the combination with uncorrelated soft particles when they are reconstructed. The existence of such kinds of jets within those containing a D meson will be discussed later in sections 5.4 and 5.7.3.

In this section, the technique for removing the average background contribution will be discussed. It consists in performing the reconstruction again using the k_t algorithm.

The k_t algorithm is very similar to the anti- k_t , but instead of using $1/p_T$ in the variables, it uses p_T , as a consequence, the algorithm starts from soft seeds and it is very adaptive to soft particles. The shape of the jets is not round in the $\eta - \phi$ phase space.

The jets obtained with the k_t algorithm are then used to calculate the average jet transverse momentum density ρ on a event-by-event basis

$$\rho = \text{median} \left\{ \frac{p_T^{\text{jet}}}{A_{\text{jet}}} \right\}, \quad (5.10)$$

where A_{jet} is the jet area in the angular $\eta - \phi$ phase space.

Usually, the two leading jets, i.e. the two jets with highest p_T , are removed from the computation of ρ . The reason is that, since these jets are likely signal jets, the background p_T density is better represented by excluding signal jets.

This momentum density is used to correct the raw-signal anti- k_t jets

$$p_T^{\text{corr}} = p_T^{\text{raw}} - \rho A_{\text{jet}}. \quad (5.11)$$

Figure 5.3: Inclusive jet transverse momentum distribution from the 0-20% most central Pb-Pb events at $\sqrt{s_{\text{NN}}} = 5.02$ TeV. The distribution is background subtracted.

The cost of removing the average background density from jets is the inserting of a *jet transversal momentum fluctuation*. The effects of the jet background fluctuation can be easily seen in the inclusive charged jet p_T distribution presented in figure 5.3 where the jet transverse momentum was corrected using equation 5.11. There are now several entries below zero and the physical interpretation is that these entries are related to jets with a transversal momentum density lower than the average p_T density in the event. These jets are not necessarily background, they can just be part of the low momentum region of the spectrum.

Jet background fluctuation can be quantified using the *Random Cones* method. This is a data-driven method where a random direction is chosen in a Pb-Pb event and a cone of radius R (the same resolution parameter used in the anti- k_T algorithm) is defined with the requirement that the center of the cone obeys $\eta^{cone} \leq 0.9 - R$ to be sure that the cone is fully within acceptance.

The transverse momentum of the cone p_T^{cone} is the sum of all particles within the cone and the distribution to quantify the background fluctuation is defined as

$$\delta p_T = p_T^{cone} - \rho \pi R^2 \quad (5.12)$$

In figure 5.4 the distribution of δp_T is presented and a gaussian fit is performed from the left side of the spectrum in order to obtain the standard deviation of the background fluctuation. This deviation is commonly used to determine the jet p_T cut. The exact value for the cut depends on the analysis and other requirements on jets. Usually three standard deviations are used. Below this value it is difficult to disentangle signal from background jets in a inclusive jet analysis, which is not the case in this thesis, since only jets tagged by D meson will be analysed.

The impact of jet background on jets containing D mesons will be better discussed later in sections 5.4 and 5.7. In the latter, the correction to background fluctuation will also be presented.

5.4 D-Tagging

Charm quark fragmentation will produce not only the D meson but also other light hadrons. The charm jet produced in the heavy-quark fragmentation has to be identified out of all those jets clusterized by FastJet.

Experiments have been using the angular correlation method, where the heavy-flavour jet is selected based on the distance in the $\eta - \phi$ phase space to a heavy meson. Another method is based on the displacement of the primary and secondary vertex of the constituents of the jet, since it is expected that some particles of the heavy-flavour jets will be displaced from the primary vertex due to their relatively long decay length.

This analysis will apply a different method for selecting jets coming from heavy quark fragmentation. The jet will be reconstructed by FastJet with a list of tracks where the two daughter tracks of the D-meson candidate are removed and a "virtual" track with the D^0 candidate kinematics is included. The method is pictured in figure 5.5.

With the applied method, the D meson is part of the jet. However, one has to be careful with the jet reconstruction because, in one single event, there is a possibility that two D-meson candidates have one track in common. For example, the same kaon in an event can be combined with two different pions in such a way that both combinations form candidates that pass cuts.

Figure 5.4: The δp_T distributions provides the magnitude of the jet background fluctuation estimated with random cones method.

Figure 5.5: Representation of the method applied to reconstruct heavy-flavour jets. The daughter tracks are replaced by the D^0 candidate itself.

In other words, there could be two D-meson candidates in the list of particles used by Fast-Jet made of three daughter particles, which is unrealistic. To avoid this inconsistency, jets are reconstructed many times for a single event, considering only one candidate at a time.

Jets selected by containing a D meson as one of their constituents are called *D-tagged jets* and will be used in this analysis. A D-tagged jet was likely originated in a process related to the hard scattering because it also produced a heavy quark, so in principle, this is by definition a signal jet. The jet background rejection in this work will be further discussed in section 5.7.3.

5.5 Signal Extraction

The raw yield of jets tagged by D mesons is performed via invariant mass analysis. The program developed for the tagging was written in a way to produce a file containing a multidimensional histogram with the raw information of the invariant mass of the D-meson candidates, their transverse momentum, the p_T of the tagged jets and the momentum fraction $z_{||}$ of the D meson with respect to the tagged jet.

Two methods were studied to extract the raw yields of the D-tagged jet transverse momentum and will be now discussed with their advantages and drawbacks. In both methods the invariant mass distribution is fit by a combination of a gaussian function for the signal and an exponential for the background

$$f(M) = f_{sig}(M) + f_{bkg}(M) = \frac{Y_{raw}}{\sqrt{2\pi}\sigma} \exp\left(-\frac{(M - m_{D^0})^2}{2\sigma^2}\right) + c_1 \exp(-c_2 M). \quad (5.13)$$

The fit parameter Y_{raw} in equation 5.13 is the extracted raw yield and c_1 and c_2 are fit parameters.

5.5.1 D-Tagged Jet p_T Direct Extraction

The invariant mass analysis is performed in bins of jet p_T with a fit function (equation 5.13). The signal extracted from the fit corresponds to the raw yield of jets tagged by a D^0 in an interval of jet transversal momentum.

The D^0 invariant mass distribution for jets with transverse momenta between -10 and 35 GeV/c is presented in figure 5.6. The D^0 candidates are limited to the range of $3 < p_T^{D^0} < 20$ GeV/c. The reason of the kinematic range is related to the efficiency corrections and will be discussed in section 5.6. All the intervals of jet p_T contain a D-meson peak with significance higher than 3.

5.5.2 Side-Band Subtraction

The invariant mass analysis is performed in bins of D-meson p_T with a fit function (equation 5.13). The jet p_T distribution related to the signal region ($\pm 3\sigma$ from the mean) of the invariant mass distribution is then taken from the file. This jet transversal momentum is proportional to the signal and the background in the signal region (red area in figure 5.7). In order to remove the background contribution, the jet transversal momentum spectrum corresponding to the side-band regions (blue area in figure 5.7) $M > 4\sigma$ and $M < -4\sigma$ from the mean are normalized to the background value extracted from the fit.

Figure 5.7 contains the D^0 invariant mass distributions and the fits, which presented significance higher than 3 in all bins. The signal region is filled in red and side-band region in blue. The jet p_T distributions are presented in figure 5.8. In red is the distribution related to the signal area in the D^0 invariant mass distribution and the blue points are the same distribution related to the normalized side-band region.

Figure 5.6: Invariant mass distributions in intervals of jet p_T from -10 to 35 GeV/c . D -meson transverse momenta are in the range of 3 to 20 GeV/c .

The normalized jet p_T distribution from the side-band region is then subtracted from the signal region spectrum, resulting in jet p_T distribution proportional to the D^0 raw yield in the considered D -meson p_T interval. The results can be seen in figure 5.9. After summing all the six spectra to obtain the D -tagged jet transverse momentum spectrum, the results can be compared to the previous method.

Figure 5.7: Invariant mass distributions in intervals of jet p_T .

5.5.3 Method Comparison

After extracting the D-tagged jet transverse momentum spectrum in both methods it is necessary to compare the results in order to test the consistency of the applied techniques.

The jet p_T distributions can be seen in figure 5.10. The spectrum obtained with the direct extraction is presented with the red points and in blue is the corresponding distribution for the side-band method. Both are compatible within less than one standard deviation.

Side-band method presents a larger statistical uncertainty due to the fact that the subtraction procedure takes into account the uncertainties of the distributions for the signal region and the side-band region.

The main differences between the methods will emerge in the procedures to correct the data by D-meson and jet reconstruction efficiency that will be discussed in the next section.

5.6 Efficiencies Corrections

Measurements can include effects from detector resolution and efficiencies. In this section the efficiencies corrections from D meson and jet reconstructions will be addressed.

Figure 5.8: Jet transverse momentum spectra for different intervals of $D^0 p_T$. The red points corresponds to the signal region and the blue points to the normalized side-band region.

Figure 5.9: Jet transverse momentum spectra for different intervals of $D^0 p_T$.

Figure 5.10: *D*-tagged jets transverse momentum. Jet p_T direct extraction method are the red points and the blue is the distribution using the side-band method.

D-meson and jet efficiencies are computed using simulations. PYTHIA 6.4 [39] was used as event generator for pp collisions and HIJING [63] was employed to Pb-Pb collision simulations. Both used GEANT3 [64] transport framework to simulate the ALICE detector response and conditions. The reconstruction algorithms are the same employed in data and, therefore, this sample will be called *reconstructed level*. Particles without reconstruction will be referred as *generated level*.

5.6.1 D-Meson Reconstruction Efficiency

The D-meson reconstruction efficiency ϵ_D is evaluated as

$$\epsilon_D(p_T) = \frac{N_{rec}^D(p_T)}{N_{gen}^D(p_T)} \quad (5.14)$$

where N_{rec}^D and N_{gen}^D are the number of reconstructed and generated D-mesons in a sample, respectively.

The reconstruction efficiency is applied in different ways depending on the signal extraction

Figure 5.11: Prompt D meson reconstruction efficiency in Pb-Pb collisions at $\sqrt{s} = 5.02$ TeV. The statistical uncertainties are smaller than the points.

method. In the case of the direct subtraction discussed in section 5.5.1, the invariant mass distribution needs to be sliced not only in intervals of jet p_T , but also in intervals of D-meson transverse momentum and then weighted by the inverse of the reconstruction efficiency and summed up back in D-meson p_T ,

$$M(p_T^{ch,jet})_{corr} = \sum_{p_T^D} \frac{M(p_T^{ch,jet}, p_T^D)}{\epsilon(p_T^D)}, \quad (5.15)$$

where M is the D-meson invariant mass distribution and $\epsilon(p_T^D)$ is the reconstruction efficiency for a given D-meson p_T .

The efficiency of prompt D mesons in Pb-Pb collision at $\sqrt{s_{NN}} = 5.02$ TeV is presented in figure 5.11. The correction for the side-band method is more straightforward since the binning is done only in D-meson p_T . The invariant mass distribution is weighted by one over the efficiency $\epsilon(p_T^D)$.

After corrections the methods are again compared as presented in figure 5.12. The points are all compatible within one sigma. The relative statistical uncertainty is also compared at this point and it is presented in figure 5.13. Even after the efficiency correction, the statistical uncertainties are still larger in the side-band method than in the direct extraction method.

Figure 5.12: *D*-tagged jet transverse momentum for both methods. The red distribution was obtained using direct extraction. The side-band method is pictured with blue points.

5.7 Unfolding

Measurements are distorted by the finite resolution of detectors and other experimental limitations, e.g. track reconstruction efficiency. Usually, these measurements are not directly comparable to theory due to these distortions.

In order to correct the distortions in a measured observable a deconvolution technique was employed. In this technique a measured distribution $\mathcal{F}(x_{\text{meas}})$ is described as a convolution of a function $\mathcal{R}(x_{\text{meas}}, x_{\text{true}})$, which contains the distortion information and a true distribution $\mathcal{T}(x_{\text{true}})$ formally defined as

$$\mathcal{M}(x_{\text{meas}}) = \int \mathcal{R}(x_{\text{meas}}, x_{\text{true}}) \mathcal{T}(x_{\text{true}}) dx_{\text{true}}. \quad (5.16)$$

Since measurements are naturally discrete, this integral equation has to be discretized in order to *unfold* the true distribution. The function $\mathcal{R}$ is treated as a matrix leading to the linear system given by

$$M = R \cdot T, \quad (5.17)$$

where M and T are vectors representing the measured and true distributions. In high-energy physics

Figure 5.13: Relative statistical uncertainties of D -tagged jet p_T for the two methods.

experiments, R is commonly called *response matrix*.

Unfortunately, the system of linear equations that results from 5.16 is an *ill-posed* problem due to the inversion of M , i.e. the result wildly oscillates due to uncertainties in the measured distribution and has no physical meaning.

There are methods to overcome the oscillating results. This thesis will discuss *unfolding* using the bayesian [65] and SVD [66] methods. The framework used is the RooUnfold [67].

5.7.1 Bayesian Method

This method is based in the Bayes' theorem. In order to express the ideas of the theorem applied to the unfolding problem, effects E_i and causes C_i are defined. For instance, a effect E_1 could be caused by C_1 or C_2 and, on the other hand, C_1 can produce the effects E_1 or E_2 . So, for an isolated effect E , the probability of this effect being produced by an specific cause C is written as $P(C, E)$.

This probability $P(C, E)$ is then proportional to the cause probability $P(C)$ to happen times the probability $P(E, C)$ of this cause to produce the effect E . So $P(C, E)$ is defined as

$$P(C, E) = \frac{P(E, C)P(C)}{\sum_l^n P(E, C_l)P(C_l)} \quad (5.18)$$

where the denominator is the total probability of all n_C possible causes C_l to produce the effect E .

It is worth noting the difference between

$P(C, E) \rightarrow$ probability of the effect E be produced by a cause C

and

$P(E, C) \rightarrow$ probability of the cause C to produce an effect E .

Usually, unfolding methods requires information *a priori* about the true distribution (*prior*) in order to regularize the process and avoid oscillating results. In the Bayesian method, this information is represented by $P(C)$ and can be somewhat vague, even starting with a uniform distribution $P_0(C)$. The knowledge of the prior increases with the number of observed events.

Considering equation 5.18 in a more general way with n_E effects and n_C causes, it is possible to write the number $n(C_i)$ of true events due to the cause C_i in a given distribution as

$$n(C_i) = \sum_{j=1}^{n_E} \frac{P(E_j, C_i)P_0(C_i)}{\sum_{l=0}^{n_C} P(E_j, C_l)P_0(C_l)} n(E_j) = \sum_{j=1}^{n_E} P(C_i, E_j) n(E_j) = \sum_{j=1}^{n_E} M_{ij} n(E_j) \quad (5.19)$$

where $n(E_j)$ is the number of measured (distorted) events E_j .

Equation 5.19 contains the elements M_{ij} of the response matrix. This matrix M is not, mathematically speaking, the inverse of the matrix R from equation 5.17.

The method works in an iterative way. First, $P_0(C_i)$ is chosen using the *a priori* information of the true distribution, though if it is unknown, an uniform distribution can be applied. This automatically defines the initial expected number of true events $n_0(C_i) = P_0(C_i)N_{obs}$, where N_{obs} is the total number of measured events. With this prior, equation 5.19 is used to calculate $n(C_i)$ and $P(C_i)$, which are used as new prior and starting the iteration again. The stopping condition can be done by analysing the χ^2 distributions or, if the prior is reasonable known, a few number of iterations are enough for convergence.

5.7.2 SVD Method

SVD stands for *Single Value Decomposition* and consists in the factorization of a $m \times n$ matrix A in

$$A = USV^T \quad (5.20)$$

with

$$UU^T = U^T U = I \quad \text{and} \quad VV^T = V^T V = I. \quad (5.21)$$

The matrix S has the elements of the main diagonal $s_{ii} \geq 0$ and those out the this diagonal are zero. The elements in the main diagonal are called singular values of the matrix A .

The problem to be solved was presented in equation 5.17 and the matrices U and V are used to rotate both sides of the equation, so the problem is decomposed as

$$M = USV^T T \quad (5.22)$$

and can be formally solved by evaluating

$$Z = V^T T, \quad Y = U^T M, \quad SZ = Y \quad \text{and} \quad Z = S^{-1} Y \quad (5.23)$$

resulting in the true spectrum written as

$$T = VZ = VS^{-1}Y = VS^{-1}U^T M = A^{-1}M, \quad (5.24)$$

and the inverse of the matrix A is $VS^{-1}U^T$.

The oscillation of the results can be put under control by choosing a prior.

5.7.3 Challenges Concerning Unfolding

Before removing smearing effects from the detectors and jet background fluctuation introduced with the background subtraction, some careful checks on the distribution to be unfolded have to be done.

More specifically, one should not unfold background. A distribution such as jet transverse momentum with a large contamination of combinatorial jets would lead to an unfolded spectrum that does not represent a real jet distribution and could not be compared to theory.

Jet analyses employ different strategies to suppress combinatorial backgrounds. One simple example is the requirement on the jet p_T to be higher than a certain value. This lower limit depends on the jet resolution parameter, since the magnitude of the jet background fluctuation is also dependent of R .

Another widely used and more sophisticated technique is the requirement of the jet to contain a constituent whose transverse momentum is higher than a minimum value. This is often used in combination with the simple jet p_T cut.

These requirements can suppress background but also remove true jets that, due to the background subtraction and fluctuation, were reconstructed in the low p_T region or even in the negative part of the spectrum. As already discussed in section 5.4, D-tagged jets are likely, by definition, non-background since the D-meson signal extraction can determine the number of signal D-mesons in a certain interval of jet p_T , completely removing the influence of background D mesons.

However, there is still another source of jet background that can contaminate the sample, which is the case of jets that, during the reconstruction in the heavy-ion environment, are split due to the large amount of soft uncorrelated particles (figure 5.14). In order to test whether the anti- k_t algorithm is robust against these background particles, an *embedding* using PYTHIA simulation and real Pb-Pb events was performed.

Figure 5.14: Signal jet with constituents in red is presented on the left. On the right, the jet is split into two jets due to the presence of soft uncorrelated particles pictured in blue.

Embedding consists of including tracks from one event to another. The mixture can be done with simulation and/or data. In this test, PYTHIA pp events were added to real Pb-Pb events. Two samples were used to reconstruct jets, one containing only tracks from simulation and another with the same simulation and data.

After reconstructing jets in a simulated event, the same event is included in a Pb-Pb event and then jets are again reconstructed, creating two samples of jets. Then a geometrical criterion in the angular $\eta - \varphi$ phase space was applied to connect the jet in one sample to another:

- jet_i^{sim} (simulation only) is the closest jet in $\eta - \varphi$ to jet_j^{emb} (embedding);
- jet_j^{emb} (embedding) is the closest jet in $\eta - \varphi$ to jet_i^{sim} (simulation only);
- the distance between jet_i^{sim} and jet_j^{emb} has to obey $\Delta R = \sqrt{(\eta_i - \eta_j)^2 + (\varphi_i - \varphi_j)^2} \leq 0.1$.

After the matching between jets in the two samples, a set of jet pairs is produced. It is worth noting that these pairs contain one jet produced only with simulated tracks (figure 5.15a) and another that can contain only simulated, only data or a (most likely) combination of tracks from simulation and data, since in the embedded sample there is simulation and data (figure 5.15b). This is pictured in figure 5.15.

(a) Jet reconstructed only with simulated tracks. **(b)** Jet reconstructed with simulated and data tracks.

Figure 5.15: Tracks from simulation are in red and tracks from data are in blue.

Every track in a given event has a *track ID*, this is used to check whether the tracks of the jets in the simulated-only sample are also inside the embedded-sample jet. Two variables are defined to quantify how much of the original (simulated only) jet is inside the embedded jet in order to understand whether the splitting of jets due to the uncorrelated particles is frequent and whether

it represents a large effect. The first is the number of tracks from the original jet that is also inside the embedded jet

$$N^{frac} = \frac{N_{pyt}^{emb}}{N_{pyt}^{sim}}, \quad (5.25)$$

where the convention N_{origin}^{sample} was used. The term *origin* indicates that the track is from PYTHIA (pyt) or data (dat) and *sample* indicates if it is simulation only (sim) or embedding (emb). Also, those particles from PYTHIA in the embedded sample that are not inside the simulated-only jet is rejected.

The second variable follows the same idea, but it was applied to the transverse momentum of the tracks

$$p_T^{frac} = \frac{\sum_{tracks} p_{T,pyt}^{emb}}{\sum_{tracks} p_{T,pyt}^{sim}} \quad (5.26)$$

where the summation in the numerator is over the PYTHIA tracks in the embedded sample jet that is also inside the simulated-only track checked by the track ID.

These two variables can indicate whether the jet is fairly reconstructed in a Pb-Pb environment even though it is contaminated with soft uncorrelated particles. The distribution of N^{frac} and p_T^{frac} is shown in figures 5.16 and 5.17 respectively.

Figure 5.16: Number of tracks ratio between PYTHIA tracks from the embedded and simulated-only samples.

The ratio of the number of tracks as defined above indicates that the splitting of jets is rare and the p_T fraction also gives the information that, even when not all particles are inside the jet, the loss of transversal momentum is small. In figure 5.17, 95% of the jets in the embedded sample contain more than 95% of the transversal momentum of the simulated-only jets considering only

Figure 5.17: *Transverse momentum fraction between PYTHIA tracks from the embedded and simulated-only samples.*

the tracks matched by the track ID.

The results show that anti- k_t algorithm is robust against soft uncorrelated particles, i.e. jet splitting during the reconstruction is rare and represents a small effect. The contamination from the uncorrelated particles in D-tagged jets can be corrected by unfolding and, since the general contamination from background jets is small, the low momentum region of the spectrum can be explored.

5.7.4 Response Matrix

The response matrix used to unfold the true spectrum is a combination of two effects. The first is related to the finite resolution of the detectors that smears the jet transversal momentum distribution. The other effect is the jet background fluctuation introduced by the jet background subtraction discussed in section 5.3.1.

Unfolding in one dimension requires a two dimensional response matrix. One of the axes is the reconstructed level (or smeared) axis that contains the detector effects and jet background fluctuation. The other one is the generated level (or true) axis, which is the prior to the unfolding algorithm.

The detector effects are quantified using a fully reconstructed PYTHIA pp simulation. This simulation was reconstructed using GEANT4 to emulate the detector conditions of the Pb-Pb run period. Also, the usual ALICE reconstruction algorithms for data were used.

The reason a simulation with heavy-ion collisions is not employed is due to the difficulties related to the jet background. It is necessary for the response matrix to include background fluctuations in the reconstructed level, however, in the generated level it would be difficult to separate the particles from the signal jet from the uncorrelated particles. It was verified that the track resolution is similar in pp and Pb-Pb events and, since it is simpler to reconstruct jets in a background-reduced environment, the simulation is done in proton-proton collisions.

Figure 5.18: Inclusive jet transversal momentum distributions from Pb-Pb data in blue and embedding (PYTHIA pp simulations and Pb-Pb data) in red.

In order to include the jet background fluctuation, the PYTHIA pp events were embedded in Pb-Pb events. In figure 5.18 the inclusive jet transversal momentum spectra are presented. Jets reconstructed only with data are in blue and the embedding sample is presented in red.

The smeared axis will include PYTHIA tracks and also Pb-Pb data. The true axis will contain jets clusterized only with generated level particles from PYTHIA. It was verified that the jet reconstruction resolution is different for inclusive jets and D-tagged jets. This is due to the strategy of reconstructing jets with a D-meson track instead of the daughters. The resolution is better for D-tagged jets and the response matrix has to be built only with jet containing a D meson as constituent.

Since the D-meson reconstruction efficiency is higher in PYTHIA pp simulations than in heavy-ion collisions, the response matrix is corrected by the prompt D^0 reconstruction efficiency in proton-proton collision presented in figure 5.19.

Figure 5.19: Prompt D^0 reconstruction efficiency in pp collisions at $\sqrt{s} = 5.02$ TeV.

After these corrections, the response matrix is evaluated for D-tagged jets coming from charm quarks. In simulation, the prompt D mesons can be easily separated from feed-down. The response matrix is presented in figure 5.20. The reconstructed axis, due to the jet background fluctuation, goes down to -10 GeV/c.

Figure 5.20: Response matrix for prompt D -tagged jet unfolding.

5.7.5 Unfolding Closure Tests

Before unfolding data, some closure and stability tests were performed in order to check whether the unfolding procedure is providing a reasonable result.

Closure tests were performed with two different embedded samples produced the same way the one described in the last section. The response matrix was built with one sample and the other was used as measurement. Unfolding was performed with bayesian method.

Figure 5.21: Left: true (red line) and unfolded (blue points) D -tagged jet transversal momentum distributions. Right: relative difference between true and unfolded spectra.

The applied cut on D-meson p_T was 2 GeV/c, which gives the lower limit in jet transverse momentum. The true and unfolded spectra are presented in figure 5.21, if the unfolding procedure is removing detector and jet background fluctuations effects, the unfolded spectrum should converge to the true distribution. The right side of figure 5.21 has the relative difference between the two spectra and the fluctuation is much smaller than the poissonian statistical uncertainty.

Though this test shows that the algorithm is working as expected, the spectrum used as measurement does not contain the statistical fluctuation that is expected from the signal extraction methods discussed in section 5.5 because the D mesons are selected directly from the simulation. In order to include such fluctuations each point of the measured spectrum was varied up or down according to the corresponding standard deviation obtained in the signal direct extraction (section 5.5.1). 500 spectra were produced this way and unfolded.

Figure 5.22 contains the results of 500 unfolding procedures performed with the smeared spectra. On the left there is the relative difference distribution and in the z axis the number of entries. The standard deviation of the relative difference is presented on the right and indicates that the unfolded procedure does not include significant uncertainties to those intrinsic uncertainties from the signal extraction method.

Figure 5.22: Left: relative difference between true and unfolded spectra. The z axis is the number of entries. Right: standard deviation of the relative difference distribution. The unfolding procedure was performed using 500 spectra.

The results show that the unfolding procedure is stable even in the low jet p_T region. This was a very important step in the feasibility this analysis.

5.8 Feed-Down

D mesons originate not only from charm quarks. B mesons coming from bottom quarks can decay in D mesons and other particles. These D mesons coming directly from charm fragmentation are called *prompt* and those from bottom are the *feed-down* (or simply non-prompt).

Figure 5.23: D^0 from B decays reconstruction efficiency in Pb-Pb collisions at $\sqrt{s} = 5.02$ TeV.

This analysis used only prompt D mesons. The strategy to remove the feed-down contribution from the D-meson sample is based on simulation, previous measurements and a hypothesis.

The simulation is performed using POWHEG [68] event generator. PYTHIA 6 was used for the later parton shower and fragmentation. POWHEG was chosen because it fairly reproduces FONLL calculations and experimental results [69]. In this simulation, each pp event was generated containing one $b\bar{b}$ pair. The observable of interest is the non-prompt D-tagged jet p_T cross-section. No detector reconstruction is applied at this stage.

The simulated cross-section needs to be rescaled to Pb-Pb collisions using the number of selected events collected in data and the nuclear overlap function $\langle T_{AA} \rangle$, obtained from Glauber model as mentioned in section 2.1, which is 18.8 mb^{-1} for 0-20% centrality class. Jets were clusterized using FastJet and the same strategy described in section 5.4. In order to save computational power, jets are reconstructed using any D^0 coming from B and, later, the branching ratio for the correct decay channel $D^0 \rightarrow K\pi$ was applied.

Each entry in the jet p_T spectrum was weighted by the ratio between the feed-down D-meson reconstruction efficiency in Pb-Pb collisions, presented in figure 5.23, over that for prompt D mesons presented in figure 5.11. This takes into account the fact that the feed-down D mesons in the data

sample were corrected by the prompt D meson reconstruction efficiency in Pb-Pb collisions, as presented in section 5.6.

The last ingredient that has to be added to this spectrum is the energy loss. This is specially difficult to estimate due to the lack of measurements. This work took advantage of the measurement of the non-prompt J/ψ nuclear modification factor [70]. Based on this measurement, this analysis uses the hypothesis that the R_{AA} of non-prompt D mesons is two times larger than the R_{AA} of prompt D mesons ($R_{AA}^{feed-down} = 2R_{AA}^{prompt}$). The prompt D^0 nuclear modification factor for Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV has not been published until the end of this thesis but, as presented in conferences [71], this observable is compatible within uncertainties to the same measurement done at 2.76 TeV (Figure 2.7). Therefore, the R_{AA} of prompt D mesons at 2.76 TeV was used in this work.

Figure 5.24: D^0 from B decays reconstruction efficiency in pp collisions at $\sqrt{s} = 5.02$ TeV.

In order to add detector effects and jet background fluctuations to the simulated spectrum, the non-prompt D-tagged jet spectrum was folded with a response matrix similar to that presented in section 5.7.4, but built with D mesons coming from B decays and corrected by non-prompt D-meson reconstruction efficiency.

The non-prompt D^0 reconstruction efficiency was calculated using the same simulation used to build the response matrices. The set of cuts is also the same and, as expected due to the larger decay length, the efficiency for feed-down D^0 is higher than for prompt. The feed-down D^0 reconstruction efficiency is presented in figure 5.24.

The response matrix for non-prompt D^0 is presented in figure 5.25. It is fully corrected by the reconstruction efficiency presented in figure 5.24. Not only the shape of the D^0 transversal

Figure 5.25: Response matrix built using only D^0 from B mesons.

momentum spectrum for feed-down is different from that for prompt, but also the jet p_T spectrum, therefore the necessity to build a specific response matrix.

The folding procedure is simply a matrix multiplication between the p_T spectrum of D-tagged jets from B mesons with the response matrix. Later, the spectrum is rescaled in order to have the same number of entries than before the folding procedure.

A summary of the correction is presented in equation 5.27,

$$N_c(p_{T,rec}^{ch,jet}) = N_{c+b}(p_{T,rec}^{ch,jet}) - R_b(p_{T,rec}^{ch,jet}, p_{T,gen}^{ch,jet}) \otimes \sum_{p_T^D} \frac{\varepsilon_b(p_T^D)}{\varepsilon_c(p_T^D)} R_{AA}^b(p_T^D) N_b^{POWHEG}(p_T^D, p_{T,gen}^{ch,jet}), \quad (5.27)$$

where N_c , N_b and N_{c+b} are the prompt, non-prompt and inclusive D-tagged jet p_T spectrum, R_b is the response matrix built with feed-down D mesons, ε_c and ε_b are the prompt and non-prompt D-meson reconstruction efficiencies in Pb-Pb collisions.

The fraction of non-prompt over measured D-tagged jet as a function of jet p_T is presented in figure 5.26. This result and its impact on the systematic uncertainties evaluated in section 5.9.2 will be discussed in chapter 6.

The correction was performed subtracting the non-prompt jet p_T from the measurement using the direct extraction method presented in figure 5.12. The prompt D-tagged jet p_T distribution corrected by feed-down subtraction is presented in figure 5.27 with entries divided by the bin width.

Figure 5.26: Fraction of non-prompt over measured D -tagged jets as a function of jet p_T .

Figure 5.27: Measured prompt D -tagged jet p_T distribution corrected by D -meson reconstruction efficiency and feed-down from B meson decay. The entries are divided by the bin width.

5.8.1 Unfolding Results

The measured prompt D-tagged jet p_T distribution presented in figure 5.27 was corrected by D-meson reconstruction efficiency and feed-down subtraction. The distribution is still affected by jet background fluctuations and detector finite resolution.

The response matrix presented in figure 5.20 was used to correct the measured spectrum. The unfolding method applied was the Bayesian unfolding method with regularization parameter 3, which is the same parameter applied in section 5.7.5 and presented stable performance.

The unfolding results are presented in figure 5.28. The points are the results of the bayesian unfolding procedure. In the next chapter, this result will be compared to simulations.

Figure 5.28: *Unfolded prompt D-tagged jet p_T spectrum.*

As a sanity check, the unfolding results are folded with the response matrix and compared with the original measured distribution that was unfolded. This is called *refolding* and is presented in figure 5.29. The refolded and measured distributions are compatible within uncertainties. Large deviations of the refolded with respect to the original distribution could indicate that the unfolding procedure was not successfully done.

Figure 5.29: *Refolded (blue line) and measured (blue points) distributions.*

5.9 Systematic Uncertainties

The method for the measurement of the D-tagged jet p_T presented in this work contains systematic uncertainties. This analysis identified the main sources of systematic uncertainties:

- Signal extraction
- Feed-down correction
- Unfolding

Each source of systematic uncertainty will be discussed in the following subsections.

5.9.1 Multi-Trial Signal Extraction

The D-meson signal extraction can introduce systematic uncertainties and, in order to estimate its magnitude, a multi-trial procedure was applied. This is a standard procedure in heavy-flavour analyses in ALICE. The procedure consists in vary one parameter, fit function, invariant mass edges, or strategy in the signal extraction procedure at a time and compare it with the central values. Each variation is called *trial*.

The list of variations are the following:

- Parameter variation
 - fixed σ and free mean
 - fix σ ($1.15 \times \sigma$) and free mean
 - fix σ ($0.85 \times \sigma$) and free mean
 - free σ and mean
 - free σ and fixed mean
 - fixed σ and mean
- Background function variation
 - exponential
 - second order polynomial
- Invariant mass edge variation
 - Invariant mass lower edge: 1.72, 1.74 (GeV/c²)
 - Invariant mass higher edge: 2.00, 2.03 (GeV/c²)
- Strategy variation: Bin counting
 - $3.5 \times \sigma$ around the peak
 - $4 \times \sigma$ around the peak

Only those trials with reduced χ^2 less or equal 3 are considered and the parameter variations are performed once per background fit function. At the end, the systematic uncertainty is taken as the RMS of all trials. Figure 5.30 presents the relative systematic uncertainties of the D-tagged jet signal extraction using the multi-trial method as a function of jet p_T .

Figure 5.30: Relative systematic uncertainties of the D-tagged jet signal extraction using the direct extraction and side-band methods.

5.9.2 Feed-Down

Some of the POWHEG simulation parameters were varied in order to evaluate the systematic uncertainty of the feed-down correction. The central values of the simulation were obtained with the bottom mass $m_b = 4.75 \text{ GeV}/c^2$, charm mass $m_c = 1.5 \text{ GeV}/c^2$, PDF = CT10NLO, which is a standard PDF for hadron collision physics [72], and $\mu_f = \mu_r = \mu_0 = \sqrt{m^2 + p_T^2}$, which are the factorization and renormalization parameters, respectively.

Only one variation at a time was applied. The bottom mass was varied in the range $4.5 < m_b < 5 \text{ GeV}/c^2$. The PDFs MSTW2008nnlo68cl and CT10 were also used to estimate the variation produced in the simulation due to different parton distribution functions. The factorization and renormalization parameters were varied in such a way that $0.5\mu_0 < \mu_f, \mu_r < 2\mu_0$ and $0.5 < \mu_f/\mu_r < 2$.

The relative difference between the central points and each variation can be seen in figure 5.31. Figure 5.31a presents the relative difference of the generated non-prompt D-tagged jet p_T distributions and, on the right, the same for the smeared spectra are presented in figure 5.31b. The smearing was done as discussed in section 5.8, i.e. multiplying the spectrum by the response matrix (figure 5.25).

Figure 5.31: Non-prompt D-tagged jet relative difference between the variation of the POWHEG simulation over the central points as a function of jet p_T .

The systematic uncertainty of the POWHEG simulation was taken as the maximum absolute variation from the central points among those presented in figure 5.31b in order to have a symmetric systematic uncertainty.

The R_{AA} hypothesis for the non-prompt D mesons was discussed in section 5.8. The systematic uncertainties associated to this hypothesis was evaluated by varying the $R_{AA}^{feed-down}$ in the range $R_{AA}^{prompt} < R_{AA}^{feed-down} < 3R_{AA}^{prompt}$.

The whole procedure described in section 5.8 for the feed-down was done in order to compute the variation from the central points. The results can be seen in figure 5.32, where the blue lines are the associated relative systematic uncertainty due to the non-prompt D-meson nuclear modification factor hypothesis.

Figure 5.32: *Relative systematic uncertainties of the non-prompt D-meson R_{AA} hypothesis.*

5.9.3 Unfolding

The unfolding procedure usually requires a regularization parameter, which, in bayesian unfolding, is the number of iterations and, based on the unfolding closure tests presented in subsection 5.7.5, the central values were evaluated with 3 iterations. In order to obtain the systematic uncertainty of this choice, the unfolding was also performed with 2 and 4 iterations. The results were used to evaluate the relative difference with respect to the central values.

Figure 5.33 presents the systematic uncertainty band from the unfolding procedure. It is worth noting that the upper and lower edges are asymmetric and, at the end, the largest absolute variation will be taken in order to obtain a symmetric uncertainty. Also, this uncertainty is not directly comparable to the previous uncertainties because this is an uncertainty associated to the unfolded spectrum and this is the reason why the binning (from 3 to 35 GeV/ c) is different from the previous systematic uncertainties (from -10 to 35 GeV/ c).

In the next subsection, the previous systematic uncertainties will be transported to the unfolded spectrum and all the systematic uncertainties will be combined in quadrature.

Figure 5.33: *Relative systematic uncertainties from Unfolding.*

5.9.4 Final Systematic Uncertainties

The systematic uncertainties from signal extraction, feed-down correction and R_{AA} hypothesis are combined bin by bin in quadrature as in equation 5.28

$$(\sigma_{sys}^*)^2 = (\sigma_{signal}^*)^2 + (\sigma_{FD}^*)^2 + (\sigma_{R_{AA}}^*)^2 \quad (5.28)$$

where σ_{signal}^* , σ_{FD}^* and $\sigma_{R_{AA}}^*$ are the systematic uncertainty from signal extraction, feed-down and non-prompt D-meson R_{AA} hypothesis.

This combined systematic uncertainty is then used to create two uncertainty spectra. The upper spectrum is the central points multiplied by $1+\sigma_{sys}^*$ and the lower spectrum is the central points multiplied by $1-\sigma_{sys}^*$.

These two uncertainty spectra are unfolded using the same procedure applied to the central points. The transported systematic uncertainty is calculated as the relative difference between the uncertainty spectra and the unfolded results of the central points. This uncertainty is then combined in quadrature with the unfolding systematic uncertainty. The final combined systematic uncertainty σ_{sys} is presented in figure 5.34.

Figure 5.34: Final combined relative systematic uncertainties.

Chapter 6

Conclusions

“Reality is that which, when you stop believing in it, doesn’t go away.”

— Philip K. Dick

This final chapter presents the conclusions of the measurement, considerations about the method, critics and future perspectives.

6.1 The Measurement

The unfolded spectrum has to be normalized in order to obtain the invariant yield

$$\frac{d^2N}{dp_T d\eta_{jet}} = \frac{1}{N_{events}} \frac{1}{\Delta\eta_{jet}} \frac{1}{BR} \frac{N_{jet}(p_T^{ch,jet})}{\Delta p_T^{ch,jet}}, \quad (6.1)$$

where N_{events} is the number of events, $\Delta\eta_{jet}$ is the jet pseudorapidity interval, BR is the branching ratio of the D^0 hadronic decay used in the reconstruction, $\Delta p_T^{ch,jet}$ is the transverse momentum bin width and N_{jet} is the number of prompt D^0 -tagged jets.

The D^0 -tagged jet transverse momentum invariant yield is presented in figure 6.1. The red points are the central values and the red lines are the statistical uncertainties. The red boxes are the systematic uncertainties associated to the D-meson signal extraction, POWHEG simulation and R_{AA} hypothesis for the feed-down correction and unfolding.

The unfolding procedure corrected the measurement only by the detector and jet background fluctuation effects included in the response matrix. Therefore the physics of other phenomena, e.g. energy loss, are still present in the spectrum. The depletion in high p_T when compared to the true spectrum from simulation is compatible to energy loss due to the creation of a medium. Nevertheless, in order to make any statement with respect to energy loss, this measurement has to be compared to the same analysis performed in other collision systems.

By comparing this measurement with the same analysis done in pp collisions, one could extract information related to hot nuclear matter effects and signatures of the Quark-Gluon Plasma calculating the nuclear modification factor for D-tagged jets. Cold nuclear matter effects can be studied using measurements in the pPb collision system.

These analyses in pp and pPb are ongoing and are a priority inside ALICE Collaboration. For now, a comparison of the measured D-tagged jet transverse momentum in Pb-Pb collisions at $\sqrt{s_{AA}} = 5.02$ TeV with simulation is provided in this chapter.

The simulation was performed using POWHEG and PYTHIA 6 with the same parameters used in the simulation of the feed-down correction. The only difference is that, for this comparison, the pp events at $\sqrt{s} = 5.02$ TeV were generated with a $c\bar{c}$ pair.

In order to make the comparison with data, the simulated D-tagged jet cross-section as a function of jet p_T was multiplied by the nuclear overlap function $\langle T_{AA} \rangle$ for the 0-20% centrality, which corresponds to 18.8 mb^{-1} . The systematic uncertainty was evaluated by varying the parameters of the simulation. The maximum and minimum variations with respect to the central values corresponds to the parameters (μ_f, μ_r) equal $(1, 0.5)$ and $(0.5, 1)$, respectively.

The simulation and data are compared in figure 6.1. The blue boxes are the systematic uncertainties of the simulation.

Figure 6.1: Invariant yield of D^0 -tagged jet transverse momentum spectrum. The red points are data and the red boxes are the systematic uncertainties. The simulation is presented in blue points with uncertainties represented by blue boxes.

The ratio between data and simulation was taken in order to evaluate the nuclear modification factor and it is shown in figure 6.2. The points below unity indicate the presence of hot nuclear matter effects, though the uncertainties from the simulation are large. The measurement of a baseline for direct comparison and R_{AA} calculation is still necessary for a decisive statement.

This is the first measurement of jets from charm quarks in this transverse momentum range in heavy-ion collisions. This is the first step towards the measurement of the fragmentation function of

Figure 6.2: Nuclear modification factor of D-tagged jets in function of the jet transverse momentum. Red points are the measurement over the simulated baseline. The red boxes are the systematic uncertainties from the measurement and the blue boxes represent systematic uncertainties from the simulation.

D mesons. It is important to note that no measurement of the fragmentation function in heavy-ion collisions was previously made until the end of this thesis.

6.2 Method Development

The method of measuring heavy-flavour jets in heavy-ion collisions has been developed in this thesis. Tagging jets with D mesons is an efficient tool to reject background jets, allowing the exploration of lower transverse momentum regions.

This analysis would benefit of more statistics, specifically for the p_T -differential fragmentation function studies. It is also worth noting that for most of the bins in the measurement presented in this thesis, the systematic uncertainties are higher or, at least, of the same magnitude of the statistical uncertainties.

The largest source of systematic uncertainties is related to the feed-down correction. Non-prompt D meson R_{AA} was not yet measured and it is necessary to make use of the best of the current knowledge, which is the non-prompt J/ψ measurement done by CMS collaboration. It introduces a significant amount of systematic uncertainty due to the wide uncertainty range that the non-prompt R_{AA} has to cover. The direct measurement of the non-prompt D-meson nuclear modification factor could significantly decrease the systematic uncertainties of the method presented in this work.

The R_{AA} hypothesis for non-prompt D mesons used in this thesis also considers that the jet and the D meson lose the same relative amount of energy in the interaction with the QGP. This is a reasonable assumption if the fragmentation occurs outside the QGP and the charm-quark fragmentation is hard ($p_T^{charm} \approx p_T^D$).

Using the assumption that the factorization of the processes involved in charm-quark production, energy loss, fragmentation and hadronization is valid, one can apply the nuclear modification factor of prompt D mesons to the simulation described in the previous section and compare it to data.

The results can be seen in figure 6.3. Data and simulation are compatible within uncertainties. Above 20 GeV/c, data and simulation start to deviate, though they are still compatible within two uncertainties. Larger deviations in this comparison could indicate that the factorization is not reliable and fragmentation can occur inside the hot medium. Another explanation would be that the fragmentation is not well described by PYTHIA 6.

Figure 6.3: Comparison between data and simulation with D-meson R_{AA} . The red points are data and the red boxes are the systematic uncertainties. The simulation is presented in blue points with uncertainties represented by blue boxes.

6.3 Final Considerations

This thesis presented the measurement of the D-tagged jet transverse momentum in Pb-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. The work developed in this thesis provides the method to measure heavy-flavour jets tagged by heavy-flavour hadrons. In principle, any heavy-flavour hadron, with enough statistics, can be used.

Nuclear modification factor is an interesting tool to quantify energy loss and has been widely used in heavy-ion physics. However, as discussed in this thesis, the fragmentation function is necessary to complement the information provided by this measurement. The method described in this work can be used for the measurement of the fragmentation function with some additional steps.

The combined information of the nuclear modification factor and fragmentation function can possibly solve the inconsistency in the light- and heavy-flavour hadron R_{AA} , since both are compatible when it would be expected a mass dependence.

Many aspects of the fragmentation is still a debate. There is no striking evidence that the fragmentation occurs outside the medium. If this is true, no modification due to hot nuclear matter effects in the fragmentation function is expected and, consequently, no modification in the measurement of the momentum fraction should be seen.

Other effects such as *recombination* [73, 74, 75, 76], could be studied by observing deviations in the nuclear modification factor and fragmentation function due to the lack of the jet in which the heavy hadron should be found since the charm quark would not fragment.

Recent measurements of high momentum jets in pp collisions by ATLAS Collaboration are not compatible with D mesons carrying large amounts of the jet momentum [77], which is a hint that the fragmentation function could depend on quark momentum and studies in different p_T regions should be carried out. In this aspect, ALICE can provide measurements in the low momentum region.

Appendix A

Conventions and Notations

Some common notations used in heavy-ion physics will be listed in this appendix. Also, some conventions about the geometry in ALICE experiment.

Figure A.1 has the convention of the geometry in ALICE experiment. The beam goes along the z axis and the transverse plane is formed by the x and y axes.

Figure A.1: *Geometry convention in ALICE experiment.*

- Rapidity (y)

The rapidity of a particle is defined as $y = \frac{1}{2} \ln \left(\frac{E+p_z}{E-p_z} \right)$, where E and p_z are the energy and the z component of the momentum of the particle.

- Pseudorapidity (η)

Defined as $\eta = -\ln \left[\tan \left(\frac{\theta}{2} \right) \right]$. This is the variable commonly used instead of the θ angle. There is a clear relationship between η and θ . For $\theta = 0^\circ$, which is the particle in the direction of the beam, $\eta = \infty$ and for $\theta = 90^\circ$, which is the particle in some direction in the azimuthal plane, $\eta = 0$.

- Transverse momentum (p_T)

The quadratic sum of the x and y components of the particle momentum $p_T = \sqrt{p_x^2 + p_y^2}$.

- Transverse mass (m_T)

The quadratic sum of the particle transverse momentum and mass $m_T = \sqrt{p_T^2 + m^2}$.

Bibliography

- [1] E. D. Bloom et al. High-Energy Inelastic e-p Scattering at 6° and 10° . *Phys. Rev. Lett.*, 29:930, 1969. 1
- [2] M. Breidenbach et al. Observed Behavior of Highly inelastic Electron-Proton Scattering. *Phys. Rev. Lett.*, 23:935, 1969. 1
- [3] W. G. Seligman et al. Improved Determination of α_S From Neutrino-Nucleon Scattering. *Phys. Rev. Lett.*, 79:1213, 1997. 2
- [4] P. S. Auchincloss et al. Measurement of the inclusive charged-current cross section for neutrino and antineutrino scattering on isoscalar nucleons. *Z. Phys. C*, 48:411, 1990. 2
- [5] R. Blair et al. Measurement of the Rate of Increase of Neutrino Cross Sections with Energy. *Phys. Rev. Lett.*, 51:343, 1983. 2
- [6] P. Berge et al. A measurement of differential cross-sections and nucleon structure functions in charged-current neutrino interactions on iron. *Z. Phys. C*, 49:187, 1991. 2
- [7] TASSO Collaboration. Evidence for planar events in e^+e^- annihilation at high energies. *Phys. Lett. B*, 86:243, 1979. 2
- [8] Mark J Collaboration. Discovery of Three-Jet Events and a Test of Quantum Chromodynamics at PETRA. *Phys. Rev. Lett.*, 43:830, 1979. 2
- [9] PLUTO Collaboration. Evidence for gluon bremsstrahlung in e^+e^- annihilations at high energies. *Phys. Lett. B*, 86:418, 1979. 2
- [10] JADE Collaboration. Observation of planar three-jet events in e^+e^- annihilation and evidence for gluon bremsstrahlung. *Phys. Lett. B*, 91:142, 1980. 2
- [11] J. Ellis and I. Karliner. Measuring the spin of the gluon in e^+e^- annihilation. *Nucl. Phys. B*, 148:141, 1979. 3
- [12] TASSO Collaboration. Evidence for a spin-1 gluon in three-jet events. *Phys. Lett. B*, 97:453, 1980. 3
- [13] H. Fritzsch, M. Gell-Mann and H. Leutwyler. Advantages of the color octet gluon picture. *Phys. Lett. B*, 47:365, 1973. 3
- [14] David J. Gross and Frank Wilczek. Ultraviolet Behavior of Non-Abelian Gauge Theories. *Phys. Rev. Lett.*, 30:1343, 1973. 3
- [15] H. David Politzer. Reliable Perturbative Results for Strong Interactions? *Phys. Rev. Lett.*, 30:1346, 1973. 3

- [16] Wikipedia. https://commons.wikimedia.org/wiki/File:Standard_Model_of_Elementary_Particles.svg. *public domain*, 2017. 4
- [17] ATLAS collaboration and CMS collaboration. Measurements of the Higgs boson production and decay rates and constraints on its couplings from a combined ATLAS and CMS analysis of the LHC pp collision data at $\sqrt{s} = 7$ and 8 TeV. *J. High Energ. Phys.*, 8:45, 2016. 4
- [18] C. Patrignani et al. (Particle Data Group). The Review of Particle Physics. *Chin. Phys. C*, 40, 2016. 5, 45
- [19] O.V. Zenin V.V. Ezhela, S.B. Lugovsky. Hadronic Part of the Muon g-2 Estimated on the $\sigma_{tot}^{2003}(e^+e^- \rightarrow hadrons)$ Evaluated Data Compilation. *arXiv:hep-ph/0312114v2*, 2004. 6
- [20] K. G. Wilson. Confinement of quarks. *Phys. Rev. D*, 10:2445, 1974. 8
- [21] H. Satz. Colour Deconfinement in Nuclear Collisions. 2000. 8
- [22] Cheuk-Yin Wong. *Introduction to High-energy Heavy-ion Collisions*. 1994. 10
- [23] et al. A. Bazavov. Equation of state in (2+1)-flavor QCD. *Phys. Rev. D*, 90:9, 2014. 11
- [24] S. A. Chin. Transition to hot quark matter in relativistic heavy-ion collision. *Phys. Lett. B*, 70:552, 1978. 11, 12
- [25] G. Sauer. and H. Chandra. Thermal properties of nuclei. *Nucl. Phys. A*, 264:221, 1976. 12
- [26] et al. A. Bazavov. The chiral and deconfinement aspects of the QCD transition. *Phys. Rev. D*, 85:054503, 2014. 12
- [27] M. A. Stephanov. QCD phase diagram and the critical point. *Prog. Theor. Phys. Suppl.*, 153:139, 2004. 12
- [28] S. J. Sanders M. L. Miller, K. Reygers and P. Steinberg. Glauber Modeling in High Energy Nuclear Collisions. *Ann. Rev. Nucl. Part. Sci.*, 57:205, 2007. 16
- [29] H. Wang. Study of Particle Ratio Fluctuations and Charge Balance Functions at RHIC. *PhD Thesis*, <https://arxiv.org/abs/1304.2073v1>, 2013. 16
- [30] R. A. Lacey. *Winter Workshop at Squaw Valley*, 2013. 17
- [31] S. J. Sanders M. L. Miller, K. Reygers and P. Steinberg. Glauber Modeling in High Energy Nuclear Collisions. *Phys. Rev. Lett.*, 111:102301, 2013. 18
- [32] J. D. Bjorken. Energy Loss of Energetic Partons in Quark - Gluon Plasma: Possible Extinction of High p(t) Jets in Hadron - Hadron Collisions. *FERMILAB-PUB-82-059-THY*, *FERMILAB-PUB-82-059-T*, 1982. 19
- [33] D. Schiff R. Baier and B.G. Zakharov. Energy loss in perturbative QCD. *Ann.Rev.Nucl.Part.Sci.*, 50:37, 2000. 19, 20
- [34] D.E. Kharzeev Yu.L. Dokshitzer. Heavy quark colorimetry of QCD matter. *Phys. Lett. B*, 519:199, 2001. 22
- [35] B. Kämpfer R. Thomas and G. Soff. Gluon Emission of Heavy Quarks: Dead Cone Effect. *Acta Phys. Hung. A*, 22:83, 2005. 22, 23
- [36] The ALICE collaboration. Transverse momentum dependence of D-meson production in Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. *J. High Energ. Phys.*, 1603:081, 2016. 24

- [37] The STAR Collaboration. Nuclear Modification Factor of D^0 Meson in Au+Au Collisions at $\sqrt{s_{NN}} = 200$ GeV. *Nucl. Phys. A*, 931:520, 2014. 24
- [38] Bo Andersson. *The Lund Model*. Cambridge University Press, 1998. 25
- [39] Stephen Mrenna Torbjorn Sjostrand and Peter Skands. PYTHIA 6.4 Physics and Manual. *J. High Energ. Phys.*, 0605:026, 2006. 25, 55
- [40] G. Mennessier X. Artru. String model and multiproduction. *Nucl. Phys. B*, 70:93, 1974. 25
- [41] M. G. Bowler. e^+e^- Production of heavy quarks in the string model. *Z. Phys. C*, 11:169, 1981. 26
- [42] I. Schmitt C. Peterson, D. Schlatter and P. M. Zerwas. Scaling violations in inclusive e^+e^- annihilation spectra. *Phys. Rev. D*, 27:105, 1983. 26
- [43] MARK II Collaboration. D-meson production in e^+e^- annihilation at $E_{c.m.}$ between 3.8 and 6.7 GeV. *Phys. Rev. D*, 26:2190, 1982. 26
- [44] J. Knobloch et al. H. Abramowicz, J.G.H. de Groot. Experimental study of opposite-sign dimuons produced in neutrino and antineutrino interactions. *Z. Phys. C - Particles and Fields*, 15:19, 1982. 26
- [45] J. M. Yelton et al. D^{*+} Production in e^+e^- Annihilation at 29 GeV. *Phys. Rev. Lett.*, 49:430, 1982. 26
- [46] ALEPH Collaboration. Study of Charm Production in Z Decays. *Eur. Phys. J. C*, 16:597, 2000. 26
- [47] L3 Collaboration. Measurement of the probability of gluon splitting into charmed quarks in hadronic Z decays. *Phys. Lett. B*, 476:243, 2000. 26
- [48] DELPHI Collaboration. Measurement of the gluon fragmentation function and a comparison of the scaling violation in gluon and quark jets. *Eur. Phys. J. C*, 13:573, 2000. 26
- [49] ALICE Collaboration. Measurement of jet suppression in central Pb-Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. *Phys. Lett. B*, 719:1, 2015. 27, 28
- [50] F. Krauss K. C. Zapp and U. A. Wiedemann. A perturbative framework for jet quenching. *J. High Energ. Phys.*, 03:080, 2013. 27
- [51] T. Renk. The Physics probed by the P_T Dependence of the Nuclear Suppression Factor. *Phys. Rev. C*, 88:014905, 2013. 27
- [52] J. C. Collins. Hard-scattering factorization with heavy quarks: A general treatment. *Phys. Rev. D*, 58:094002, 1998. 31
- [53] E. Norrbin and T. Sjostrand. Production and hadronization of heavy quarks. *Eur. Phys. J.*, C17:137–161, 2000. 31
- [54] N. Armesto. Nuclear shadowing. *J. Phys. G*, 32:R367, 2006. 32
- [55] Michele Arneodo. Nuclear effects in structure functions. *Physics Reports*, 240:301, 1994. 33
- [56] ALICE Collaboration. D-meson production in p-Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV and in pp collisions at $\sqrt{s} = 7$ TeV. *Phys. Rev. C*, 94:054908, 2016. 33

- [57] CERN. <http://te-epc-lpc.web.cern.ch/te-epc-lpc/machines/pagesources/Cern-Accelerator-Complex.jpg>. *CERN*, 2013. 35
- [58] ALICE Collaboration. ALICE: Physics Performance Report, Volume I. *J. Phys. G: Nucl. Part. Phys.*, 30:1517, 2004. 36, 37
- [59] R. Schicker. Overview of ALICE results in pp, pA and AA collisions. *EPJ Web Conf.*, 138:01021, 2017. 40, 41
- [60] Pierre Billoir. Track fitting with multiple scattering: A new method. *Nucl. Instrum. Meth.*, 225:352, 1984. 40
- [61] E. Abbas et al. Performance of the ALICE VZERO system. *JINST*, 8:P10016, 2013. 41, 42
- [62] Matteo Cacciari, Gavin P. Salam, and Gregory Soyez. FastJet User Manual. *Eur. Phys. J.*, C72:1896, 2012. 46
- [63] Xin-Nian Wang and Miklos Gyulassy. hijing. *Phys. Rev. D*, 44:3501–3516, Dec 1991. 55
- [64] R. Brun, F. Bruyant, M. Maire, A. C. McPherson, and P. Zancarini. GEANT3. 1987. 55
- [65] G. D’Agostini. A Multidimensional Unfolding Method Based on Bayes’ Theorem. *Nucl. Instrum. Meth.*, 362:487, 1995. 58
- [66] Andreas Höcker and Vakhtang Kartvelishvili. SVD Approach to Data Unfolding. *Nucl. Instrum. Meth.*, 372:469, 1996. 58
- [67] Tim Adye. Unfolding algorithms and tests using RooUnfold. *arXiv:1105.1160*. 58
- [68] Simone Alioli, Paolo Nason, Carlo Oleari, and Emanuele Re. A general framework for implementing NLO calculations in shower Monte Carlo programs: the POWHEG BOX. *JHEP*, 06:043, 2010. 68
- [69] Matteo Cacciari, Stefano Frixione, Nicolas Houdeau, Michelangelo L. Mangano, Paolo Nason, and Giovanni Ridolfi. Theoretical predictions for charm and bottom production at the LHC. *JHEP*, 10:137, 2012. 68
- [70] Vardan Khachatryan et al. Suppression and azimuthal anisotropy of prompt and nonprompt J/ψ production in PbPb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV. *Eur. Phys. J.*, C77(4):252, 2017. 69
- [71] ALICE Collaboration. D-meson nuclear modification factor and elliptic flow measurements in Pb-Pb collisions at $\sqrt{s_{\text{NN}}} = 5.02$ TeV with ALICE at the LHC. *Nuclear Physics A*, 967(Supplement C):612 – 615, 2017. The 26th International Conference on Ultra-relativistic Nucleus-Nucleus Collisions: Quark Matter 2017. 69
- [72] Hung-Liang Lai, Marco Guzzi, Joey Huston, Zhao Li, Pavel M. Nadolsky, Jon Pumplin, and C. P. Yuan. New parton distributions for collider physics. *Phys. Rev.*, D82:074024, 2010. 76
- [73] Rudolph C. Hwa and C. B. Yang. Scaling behavior at high p_T and the p/π ratio. *Phys. Rev. C*, 67:034902, Mar 2003. 85
- [74] Rudolph C. Hwa and C. B. Yang. Recombination of shower partons in fragmentation processes. *Phys. Rev. C*, 70:024904, Aug 2004. 85
- [75] Rudolph C. Hwa and C. B. Yang. Recombination of shower partons at high p_T in heavy-ion collisions. *Phys. Rev. C*, 70:024905, Aug 2004. 85

- [76] Rudolph C. Hwa. Hadron Correlations in Jets and Ridges through Parton Recombination. In Rudolph C. Hwa and Xin-Nian Wang, editors, *Quark-gluon plasma 4*, pages 267–322. 2010. 85
- [77] ATLAS Collaboration. Measurement of $D^{*\pm}$ Meson Production in Jets from pp Collisions at $\sqrt{s} = 7$ TeV with the ATLAS Detector. *Phys. Rev. D*, 85:052005, 2012. 85