Deconfinement and degrees of freedom in \( pp \) and \( A - A \) collisions at LHC energies

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Abstract We present the extraction of the temperature by analyzing the charged particle transverse momentum spectra in lead–lead (Pb–Pb) and proton–proton (\( pp \)) collisions at LHC energies from the ALICE Collaboration using the Color String Percolation Model (CSPM). From the measured energy density \( \varepsilon \) and the temperature \( T \) the dimensionless quantity \( \varepsilon / T^4 \) is obtained to get the degrees of freedom (DOF), \( \varepsilon / T^4 = \text{DOF} \times \pi^2 / 30 \). We observe for the first time a two-step behavior in the increase of DOF, characteristic of deconfinement, above the hadronization temperature at temperature \( \sim 210 \text{ MeV} \) for both Pb–Pb and \( pp \) collisions and a sudden increase to the ideal gas value of \( \sim 47 \) corresponding to three quark flavors in the case of Pb–Pb collisions.

1 Introduction

The Quantum Chromodynamics (QCD) phase diagram is closely related to the history of the universe and can be probed by heavy ion collisions. Of particular interest in the heavy ion collision experiments are the details of the deconfinement and chiral transitions which determine the QCD phase diagram. One of the main challenges of the field is to simultaneously determine the temperature and the energy density of the matter produced in a collision and hence the number of thermodynamic degrees of freedom (DOF) [1]. The present work explores the initial stage of high energy collisions at LHC energies analyzing the published ALICE data [2,3] on the transverse momentum \( (p_t) \) spectra of charged particles using the framework of the clustering of color sources (CSPM) [4]. This approach has been successfully used to describe the initial stages in the soft region in high energy nucleus–nucleus and nucleon–nucleon collisions [4–9]. The CSPM is in fact different from the hydrodynamics picture and is more in line with other studies where the interaction among strings [10–12] or the domain color structure [13,14] is taken into account. The determination of the DOF requires the measurement of the initial thermalized (maximum entropy) temperature and the initial energy density at time \( \sim 1 \text{ fm/c} \) of the hot matter produced in high energy A–A and \( pp \) collisions. Lattice Quantum Chromodynamics simulations (LQCD) indicate that the non-perturbative region of hot QCD matter extends up to temperature of 400 MeV, well above the universal hadronization temperature [15].

2 Clustering of color sources and percolation

Multiparticle production is currently described in terms of color strings stretched between the projectile and the target, which decay into new strings and subsequently hadronize to produce observed hadrons. Color strings may be viewed as small areas in the transverse plane filled with color field created by colliding partons. With growing energy and size of the colliding system, the number of strings grows, and they start to overlap, forming clusters, in the transverse plane very much similar to disks in two dimensional percolation theory [16]. At a certain critical density \( \xi_c \sim 1.2 \) a macroscopic cluster appears that marks the percolation phase transition. For nuclear collisions, this density corresponds to \( \xi = N_s S_1 / S_A \) where \( N_s \) is the total number of strings created in the collision, each one of area \( S_1 = \pi r_0^2 \) and \( S_A \) corresponds to nuclear overlap area, with \( r_0 \approx 0.2 \text{ fm} \). This is the Color String Percolation Model (CSPM) [17,18].
The interaction between strings occurs forming clusters when they overlap and the general result, due to the SU(3) random summation of charges, is a reduction in multiplicity and an increase in the average transverse momentum squared, \( \langle p_T^2 \rangle \). The strings decay into new ones through color neutral \( q - \bar{q} \) pairs production. The Schwinger QED\(_2\) string breaking mechanism produces these \( q - \bar{q} \) pairs at time \( \tau_{\text{pro}} \sim 1 \) fm/c, which subsequently hadronize to produce the observed hadrons [19]. Schwinger mechanism has also been used in the decay of color flux tubes produced by the quark-gluon plasma for modeling the initial stages in heavy ion collisions [20,21].

The combination of the string density dependent cluster formation and the 2D percolation clustering phase transition, are the basic elements of the non-perturbative CSPM. The percolation theory governs the geometrical pattern of string clustering. Its observable implications, however, require the introduction of some dynamics in order to describe the behavior of the cluster formed by several overlapping clusters. We assume that a cluster of \( n \) strings that occupies an area of \( S_n \) behaves as a single color source with a higher color field \( Q_n \) corresponding to the vectorial sum of the color charges of each individual string \( Q_i \). The resulting color field covers the area of the cluster. As \( Q_n = \sum_i^n Q_i \), and the individual string colors may be oriented in an arbitrary manner respective to each other, the average \( Q_i Q_j \) is zero, and \( Q_n^2 = n Q_i^2 \).

Knowing the color charge, one can compute the multiplicity \( \mu_n \) and the mean transverse momentum squared \( \langle p_T^2 \rangle_n \) of the particles produced by a cluster, which are proportional to the color charge and color field, respectively [18]

\[
\mu_n = \sqrt{\frac{n S_n}{S_1}} \mu_0; \quad \langle p_T^2 \rangle_n = \sqrt{\frac{n S_1}{S_n}} \langle p_T^2 \rangle_1,
\]

where \( \mu_0 \) and \( \langle p_T^2 \rangle_1 \) are the mean multiplicity and transverse momentum squared of particles produced from a single string with a transverse area \( S_1 = \pi r_0^2 \) with \( r_0 = 0.2 \) fm [4]. For strings just touching each other \( S_n = n S_1 \), and \( \mu_n = n \mu_0 \), \( \langle p_T^2 \rangle_n = \langle p_T^2 \rangle_1 \). When strings fully overlap, \( S_n = S_1 \) and therefore \( \mu_n = \sqrt{n} \mu_0 \) and \( \langle p_T^2 \rangle_n = \sqrt{n} \langle p_T^2 \rangle_1 \), so that the multiplicity is maximally suppressed and the \( \langle p_T^2 \rangle_n \) is maximally enhanced. This implies a simple relation between the multiplicity and transverse momentum \( \mu_n (\langle p_T^2 \rangle_n = n \mu_0 (\langle p_T^2 \rangle_1) \), which means conservation of the total transverse momentum produced. In the thermodynamic limit, one obtains the average value of \( n S_1 / S_n \) for all the clusters [17, 18]

\[
\frac{n S_1}{S_n} = \frac{(\xi)^{1/2}}{1 - e^{-\xi}} = \frac{1}{F(\xi)^2},
\]

where \( F(\xi) \) is the color suppression factor by which the overlapping strings reduce the net color charge of the strings. With \( F(\xi) \rightarrow 1 \) as \( \xi \rightarrow 0 \) and \( F(\xi) \rightarrow 0 \) as \( \xi \rightarrow \infty \), where \( \xi = \frac{N_s S_1}{S_n} \) is the percolation density parameter. Equation (1) can be written as \( \mu_n = n F(\xi) \mu_0 \) and \( \langle p_T^2 \rangle_n = n \langle p_T^2 \rangle_1 / F(\xi) \).

It is worth noting that CSPM is a saturation model similar to the Color Glass Condensate (CGC), where \( \langle p_T^2 \rangle_1 / F(\xi) \) plays the same role as the saturation momentum scale \( Q_s^2 \) in the CGC model [22, 23].

### 3 Extraction of the color suppression factor \( F(\xi) \)

In the present work we have extracted \( F(\xi) \) in \( \text{Pb–Pb} \) collisions using ALICE data from the transverse momentum spectra of charged particles at \( \sqrt{s_{NN}} = 2.76 \) and 5.02 TeV at various centralities [3]. In case of \( pp \) collisions at \( \sqrt{s} = 13 \) TeV, \( F(\xi) \) has been obtained in high multiplicity events [2]. Only the softer sector of the spectra with \( p_t \) in the range 0.15–1.0 GeV/c is considered. We have checked that the extension of the soft range up to \( p_t = 2 \) GeV/c does not change significantly our results. To evaluate the initial value of \( F(\xi) \) from data a parameterization of the experimental data of \( p_t \) distribution in low energy \( pp \) collisions at \( \sqrt{s} = 200 \) GeV, where strings have low overlap probability, was used [5]. The charged particle spectrum is described by a power-law [4]

\[
d^2 N_c / dp_t^2 = a / (p_0 + p_t)\alpha, \quad (3)
\]

where \( a \) is the normalization factor, \( p_0 \) and \( \alpha \) are fitting parameters with \( p_0 = 1.98 \) and \( \alpha = 12.87 \) [5]. This parameterization is used both in \( \text{Pb–Pb} \), and in \( pp \) collisions to take into account the interactions of the strings [4]. The parameter \( p_0 \) in Eq. (3) is for independent strings and gets modified

\[
p_0 \rightarrow p_0 \left( \frac{\langle n S_1 / S_n \rangle_{\text{mod}}}{\langle n S_1 / S_n \rangle_{pp}} \right)^{1/4}, \quad (4)
\]

In \( pp \) collisions at low energies only two strings are exchanged with low probability of interactions, so that \( \langle n S_1 / S_n \rangle_{pp} \approx 1 \), which transforms Eq. (3) into

\[
d^2 N_c / dp_t^2 = \frac{a}{(p_0 \sqrt{1 / F(\xi)^{\text{mod}} + p_t})^\alpha}, \quad (5)
\]

where \( F(\xi)^{\text{mod}} \) is the modified color suppression factor and is used in extracting \( F(\xi) \) both in \( \text{Pb–Pb} \) and \( pp \) in high multiplicity events. The color suppression factor \( F(\xi) \) encodes the effect of the interaction among strings once they overlap. In the thermodynamic limit \( F(\xi) \) is related to the string density \( \xi \) [4]

\[
F(\xi) = \sqrt{\frac{1 - e^{-\xi}}{\xi^2}}. \quad (6)
\]

The partons produced in the decay of a cluster interact with the color of the clusters in their way out the collision surface, modifying their momenta. This reproduces the experimental harmonics of the azimuthal asymmetry observed in
pp, p - A, and A-A collisions. The size of the momentum modification is small and do not change the $p_T$ distribution.

4 Results and discussions

Figure 1a shows $F(\xi)$ as a function of $N_{\text{tracks}}/\Delta\eta$ scaled by the transverse area $S_\perp N_{\text{part}}$ is the charged particle multiplicity in the pseudorapidity range $|\eta| < 0.8$ with $\Delta\eta = 1.6$ units. For $pp$ collisions $S_\perp$ is multiplicity dependent as obtained from IP-Glasma model [28]. In case of Pb–Pb collisions the nuclear overlap area was obtained using the Glauber model [31]. A universal scaling behavior is observed in hadron-hadron and nucleus-nucleus collisions. The percolation density parameter $\xi$ is obtained from $F(\xi)$ using the relation Eq. (6) is shown in Fig. 2 as a function of number of participants $N_{\text{part}}$ for Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ and 5.02 TeV. It is observed that $\xi$ rises slowly at higher $N_{\text{part}}$. This behavior is similar to measured $(dN_{ch}/d\eta)/N_{\text{part}}$ as shown in Fig. 2.

The connection between $F(\xi)$ and the temperature $T(\xi)$ involves the Schwinger QED$_2$ mechanism (SM) for particle production. The Schwinger distribution for mass less particles is expressed in terms of $p_T^2$ [4,32]

$$dn/dp_T^2 \sim e^{-\pi p_T^2/x^2} \quad (7)$$

where the average value of the string tension is $\langle x^2 \rangle$. The tension of the macroscopic cluster fluctuates around its mean value because the chromoelectric field is not constant. The origin of the string fluctuation is related to the stochastic nature of the QCD vacuum. Since the average value of the color field strength must vanish, it cannot be constant but changes randomly from point to point [33,34]. Such fluctuations lead to a Gaussian distribution of the string tension and transforms SM into the thermal distribution [33]

$$dn/dp_T^2 \sim \exp \left( -p_T \sqrt{2\pi/\langle x^2 \rangle} \right), \quad (8)$$

with $\langle x^2 \rangle = \pi \langle p_T^2 \rangle_1/F(\xi)$. The temperature is expressed as

$$T(\xi) = \sqrt{\langle p_T^2 \rangle_1/2F(\xi)}. \quad (9)$$

We adopt the point of view that the universal hadronization temperature is a good measure of the upper end of the cross over phase transition temperature $T_h$ [35]. The single string average transverse momentum $\langle p_T^2 \rangle_1$ is calculated at $\xi = 1.2$ with the universal hadronization temperature $T_h = 167.7 \pm 2.6$ MeV [35]. This gives $\sqrt{\langle p_T^2 \rangle_1} = 207.2 \pm 3.3$ MeV.

Figure 1b shows a plot of temperature as a function of $N_{\text{tracks}}/\Delta\eta$ scaled by $S_\perp N_{\text{tracks}}$ is the charged particle multiplicity in the pseudorapidity range $|\eta| < 0.8$ with $\Delta\eta = 1.6$ units. The temperature dependence, for the systems investigated, falls on a universal curve as a function of the scaled multiplicity. The horizontal line at $\sim 165$ MeV is the universal hadronization temperature obtained from the systematic comparison of the statistical thermal model parametrization of hadron abundances measured in high energy $e^+e^-$, pp, and A-A collisions [35]. In Fig. 1b the temperatures obtained
for Pb–Pb and pp are above the hadronization temperature indicating deconfinement.

Recently, it has been suggested that fast thermalization in A–A and pp collisions can occur through the existence of an event horizon due to a rapid deceleration of the colliding nuclei [36,37]. The thermalization in this case is due to the Hawking–Unruh effect [36,38–40]. In CSPM the strong color field inside the large cluster produces deceleration of the primary $q\bar{q}$ pair which can be seen as a thermal temperature by means of Hawking–Unruh effect [39,40]. In the CSPM the initial local temperature determined at the string level becomes the global temperature of the initial state as the cluster spans the whole area above the percolation threshold.

The QGP according to CSPM is born in thermal equilibrium because the initial temperature is determined at the string level. Above the critical temperature $T > T_c$ the CSPM energy expands according to the Bjorken boost invariant 1D hydrodynamics [41]

$$\varepsilon = \frac{3}{2} \frac{dN_c}{d\tau} \langle m_f \rangle \tau_{pro}^{-1},$$  \hspace{1cm} (10)

where $\varepsilon$ is the energy density, $S_\perp$ the nuclear overlap area, $m_f$ is the transverse mass and $\tau_{pro}$ is the Schwinger production time for a boson (gluon) $\tau_{pro} = 2.45/(m_f)$ [42]. Figure 3 shows $\varepsilon$ as a function of $\xi$ for Pb–Pb and pp collisions. We observe a slow rise of $\varepsilon$ for low values of $\xi$ followed by a faster rise later. This is due to the nonlinear increase in multiplicity at higher $\xi$ values. In the case of Pb–Pb these data show a large departure of the scaling of the multiplicity per participant. This change occurs at $\xi \sim 4.2$ and $\sim 5.0$ for Pb–Pb at 2.76 TeV and 5.02 TeV respectively. For pp the jump in $\varepsilon$ is at $\xi \sim 3.8$.

In Fig. 4 we show the results of dimensionless quantity $\varepsilon/T^4$ for Pb–Pb collisions at two different energies and those for pp at 13 TeV. The results are compared with LQCD predictions for 2+1 flavors [43,44]. It is observed that CSPM results agree with LQCD results up to the temperature of $T \sim 210$ MeV for the Pb–Pb collisions. Beyond this temperature the $\varepsilon/T^4$ in CSPM rises much faster and reaches the ideal gas value of $\varepsilon/T^4 \sim 16$ at $T \sim 230$ MeV. In this region, there is a strong screening due to the large degree of overlapping of the strings, producing a faster approach to the quark gluon gas limit.

The DOF are obtained using the relation $\varepsilon/T^4 = \text{DOF} \pi^2/30$ [19]. At $T \sim 210$ MeV, $\varepsilon/T^4 \sim 11$ which corresponds to $\sim 33$ DOF while at $T \sim 230$ MeV there are $\sim 47$ DOF. It is observed that Pb–Pb at $\sqrt{s_{NN}} = 2.76$ TeV has similar features as seen at 5.02 TeV. In pp collisions at $\sqrt{s} = 13$ TeV only $\sim 33$ DOF are reached. Our results are in agreement with the conclusions obtained studying the trace anomaly in a quasi particle gluonic model [45,46]. In this model the DOF of the free gluons are also obtained for $T \simeq 1.3 T_c (T_c \simeq 165$ MeV). There are several reasons that can explain the disagreement between CSPM and lattice QCD for $T > 1.3 T_c$. First in A-A and pp collisions a strong magnetic field $B$ is produced, which is larger for A-A than for pp and increases with energy. The lattice studies of the chiral phase structure of three flavor QCD in a background magnetic field show that chiral condensate and the phase transition temperature always increase with $B$. The transition becomes stronger and turns into a first order instead of crossover [47]. As $B$ is higher in PbPb than in
pp, we expect a higher phase transition temperature in these cases.

Recently, it has been pointed out using lattice simulations that in addition of the standard crossover phase transition at $T \sim 155$ MeV, the existence of a new infrared phase transition at temperature $T$, $200 < T_{IR} < 250$ MeV. In this phase, asymptotic freedom works and therefore there is no interaction. In between these two temperatures there is coexistence of the short and long distance scales [48], which supports the present observation in our work.

A new phase in QCD also has been proposed studying the Dirac operator. While confining chromo-electric interaction is distributed among all modes of Dirac operator, the chromo-magnetic interaction is located predominantly in the near zero modes. Above $T \sim 155$ MeV the near zero modes are suppressed but not the rest of the modes, surviving the chromo-electric interaction which is suppressed at higher temperature [49].

5 Conclusion

We have used the Color String Percolation Model (CSPM) to explore the initial stage of high energy nucleus-nucleus and nucleon-nucleon collisions and determined the thermalized initial temperature of the hot nuclear matter at an initial time $\sim 1$ fm/c. For the first time the temperature and the energy density of the hot nuclear matter, from the measured charged particle spectra using ALICE data for Pb–Pb collisions, from the measured charged particle spectra using ALICE data for Pb–Pb collisions at $\sqrt{s_{NN}} = 2.76$ and 5.02 TeV and pp collisions at $\sqrt{s} = 13$ TeV, have been obtained. The dimensionless quantity $\varepsilon / T^4$ is evaluated to obtain the number of degrees of freedom (DOF) of the deconfined phase. We observe two features hitherto not reported: the existence of two temperature ranges in the behavior of the Pb–Pb system DOF, and a clear departure from the LQCD results regarding the maximum number of DOF, which reaches values in agreement with the Stephan Boltzmann limit for an ideal gas of quarks and gluons.

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