Thermal and hard scales in transverse momentum distributions, fluctuations and entanglement

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We analyze the transverse momentum distributions of pp, pPb, XeXe and PbPb collisions at different RHIC and LHC energies and centralities as well as the corresponding distributions for Higgs production decaying into γγ and 4l. A simple linear relation is found between the effective thermal temperature and the hard scale, approximately valid for all processes and mainly determined by the hard scale fluctuations. In order to go further, it is shown that the whole spectrum of pp collisions can be described by a single function showing that the thermal temperature is determined solely by the hard scale and its fluctuations. The possible relation between the multiplicities of the soft and hard scales is explored.

I. INTRODUCTION

The apparent thermal features of proton-proton collisions is a challenge to understand the collective behavior observed in small systems, where the application of the conventional hydrodynamical explanation is questionable [1, 2]. On the other hand, theoretical studies of quenches in entangled systems described by (1+1)-dimensional conformal field theories of expanding quantum fields and strings, show that the system obeys a generalized Gibbs ensemble with an effective temperature set by the energy cut-off for the ultraviolet modes [3–6]. In the last years there has been a large activity in the field [7–17].

In a high energy collision a hard parton interaction produces a rapid quench of the entangled partonic state [18] and thus the corresponding effective temperature, inferred from the exponential shape of the transverse momentum distribution (TMD) of the secondaries of the momentum distribution (TMD) of the secondaries of the hard scale. The hard process, with transverse momentum \( p_\perp \), probes only the region of the space \( H \) of transverse size \( 1/p_\perp \). Let us denote by \( S \) the region of space complementary to \( H \). The initial state is described by the wave function

\[
|\Psi\rangle = \sum_n \alpha_n |\Psi^H_n\rangle \otimes |\Psi^S_n\rangle,
\]

of a suitably chosen orthonormal set of states \( |\Psi^H_n\rangle \) and \( |\Psi^S_n\rangle \) localized in the domains \( H \) and \( S \), with different numbers \( n \) of partons. The state (1) cannot be separated into a product \( |\Psi^H\rangle \otimes |\Psi^S\rangle \), and therefore \( |\Psi\rangle \) is entangled. The density matrix of the mixed state probed in the region \( H \) is

\[
\rho_H = \text{Tr}_S \rho = \sum_n \langle \Psi^S_n | \Psi \rangle \langle \Psi | \Psi^S_n \rangle = \sum_n |\alpha_n|^2 |\Psi^H_n\rangle \langle \Psi^H_n|,
\]

where \( |\alpha_n|^2 \equiv p_n \) is the probability of having a state with \( n \) partons. We can consider that a high momentum partonic configuration of the initial state when the interaction takes place undergoes a rapid quench. The onset \( \tau \) of this hard interaction is given by the hardness scale, \( \tau \sim 1/p_\perp \). Because \( \tau \) is small the quench creates a highly excited multi-particle state. The produced particles have a thermal like exponential spectrum with an effective temperature which is determined by the hard scale and the fluctuations on the number of partons. With these considerations, we fit the different TMDs by an exponential distribution and a power like distribution [19,20,30],

\[
\frac{1}{N_{ee}} \frac{1}{2\pi p_\perp} \frac{d^2 N_{ev}}{d\eta dp_\perp} = A_{th} e^{-m_{th}/T_{th}} + \frac{A_h}{(1 + m_h^2/kT_h^2)^2},
\]

where \( T_{th} \) is the effective temperature, \( T_h \) is the hard scale and \( k \) a parameter which is determined by the falloff of the different distributions at high \( p_\perp \).
TABLE I: Temperatures and falloff $k$ of the TMD of charged particles in $pp$ collisions at different RHIC and LHC colliding energies.

| $\sqrt{s_{nn}}$ (GeV) | $T_{th}$ (GeV) | $T_h$ (GeV) | $k$  |
|------------------------|----------------|-------------|------|
| 64 GeV                 | 0.186          | 0.719       | 5.50 |
| 200 GeV                | 0.189          | 0.752       | 4.62 |
| 900 GeV                | 0.176          | 0.693       | 3.62 |
| 2.76 TeV               | 0.180          | 0.713       | 3.28 |
| 5.02 TeV               | 0.184          | 0.735       | 3.13 |
| 7 TeV                  | 0.180          | 0.716       | 3.08 |
| 13 TeV                 | 0.181          | 0.744       | 3.01 |

In the next sections, we will show that the relation between the two scales is given by the size of the fluctuations of the hard scale. This relation is universal, valid for $pp$, $pPb$, $XeXe$, and $PbPb$ for all energies and centralities. In the case of $pp$ collisions, we are able to describe the whole $p_{\perp}$ spectrum with a single function for energies ranging from RHIC to LHC. This function provides us with the low $p_{\perp}$ fluctuations needed to improve the universal agreement found in $AA$ and $pA$ collisions.

II. THE THERMAL TEMPERATURE DETERMINED BY THE HARD SCALE AND ITS FLUCTUATIONS

The results of our fit of the TMDs to [3] are shown in tables, and the corresponding uncertainties of the extracted parameters are shown as error bands in the accompanying figures. In Table [II] we show the values of $T_{th}$, $T_h$ and $k$ obtained from the fits to the TMD of charged particles produced in $pp$ collisions at very different RHIC and LHC energies. In Table [II] we show the results from the fits to the TMD of charged pions for several multiplicity bins in $pp$ collisions at $\sqrt{s_{nn}}=7$ TeV, in the range $|\eta| < 0.5$ and $0.1 < p_{\perp} < 17.5$ GeV/c. The obtained values of temperature are larger for charged particles than for pions, as expected. In Table [III] we show the fit results for $pPb$ collisions at $\sqrt{s_{nn}}=5.02$ TeV in the range $-0.5 < \eta < 0$ and $0.1 < p_{\perp} < 17.5$ GeV/c. In Table [IV] we show the fit results for $XeXe$ collisions at $\sqrt{s_{nn}}=5.44$ TeV at different centralities in the range $\eta < 0.8$ and $0.2 < p_{\perp} < 17$ GeV/c, and in Table [V] for $PbPb$ collisions at $\sqrt{s_{nn}}=5.02$ TeV at different centralities in the range $\eta < 0.8$ and $0.2 < p_{\perp} < 17$ GeV/c. The general behavior of the temperatures $T_{th}$ and $T_h$ is to increase with the colliding energy, as expected. In the case of the hard scale $T_h$, it also increases with the centrality in all cases. However, $T_{th}$ smoothly decreases with the centrality in $pp$ collisions, contrary to the rest of cases. The behavior of the falloff index $k$ with centrality is the same as the thermal temperature $T_{th}$, it increases with centrality except in $pp$ collisions, where it decreases. From these

FIG. 1: (color online) Thermal scale $T_{th}$ compared to the quenched hard scale $T_h/(k+1)$, extracted from RHIC and LHC collected data on $pp$ collisions, as a function of the collision energy.

FIG. 2: (color online) Thermal scale $T_{th}$ compared to the quenched hard scale $T_h/(k+1)$, extracted ALICE collected data on $pp$ collisions at $\sqrt{s_{nn}}=7$ TeV, as a function of the charged particle multiplicity.
values we have found the approximate relation

\[ \frac{T_h}{k + 1} = T_{th}. \]  

(4)

In Figure 1, we plot \( T_{th} \) and \( T_h/(k + 1) \) for \( pp \) collisions at different energies. In Figures 2, 3, and 4, we plot the cases of \( pp, pPb, XeXe, \) and \( PbPb \) collisions at \( \sqrt{s_{nn}} = 7 \) TeV, 5.02 TeV, and 5.44 TeV for different identified charged particles as a function of the multiplicity. We have looked at the transverse momentum distribution of Higgs production decaying into \( \gamma \gamma \) and 4\( \ell \), as reported in reference [19], obtaining \( T_{th} = 3.5 \pm 0.7 \) GeV, \( T_h = 14.4 \pm 0.3 \) and \( k = 3.7 \pm 0.4 \). Hence we obtain a quenched hard scale \( T_h/(k + 1) = 3.1 \pm 0.4 \) GeV, to compare with \( T_{th} = 3.5 \pm 0.7 \) GeV. In the same line, studies of the \( p_\perp \) distribution of the W’s produced in \( pp \) collisions have found similar behaviors [31,32].

An overall agreement is observed in all cases, except in very low multiplicity events and very low colliding energies in \( pp \) collisions, where the discrepancies between both quantities becomes substantial. This agreement is remarkable, considering the large number of TMDs studied and the large differences presented in the temperatures for different projectiles and targets, as well as centralities and energies.

In order to improve the agreement with \( pp \) data, we devise now a simple expression describing well the whole soft and hard spectrum for the full range of energies explored at RHIC and the LHC. The hard part of the TMD (3) can be rewritten as

\[ \frac{1}{1 + (p_\perp^2/\gamma)^k} = \int_0^\infty dx e^{-p_\perp^2 x} \frac{\Gamma(k)}{\Gamma(k)} (\gamma x)^{k-1} e^{-\gamma x} = \int_0^\infty dx f(p_\perp, x) W_p(x), \]  

(5)

where \( W_p(x) \) is the Gamma distribution and thus \( 1/k \) can be understood as the normalized fluctuations of the hard scale \( T_h \)

\[ \frac{1}{k} = \frac{\langle x^2 \rangle - \langle x \rangle^2}{\langle x \rangle^2}, \quad \gamma \equiv kT_h^2. \]  

(6)

We can add an additional source of fluctuations for \( T_h \) using a Gaussian distribution \( G(T_h) \). The whole \( p_\perp \) dis-

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**TABLE III:** Characteristic temperatures and falloff \( k \) of the TMD of \( \pi^\pm \) in \( pPb \) collisions at \( \sqrt{s_{nn}} = 5.02 \) TeV for several multiplicities.

| \( dN_{ch}/d\eta \) | \( T_{th}(\text{GeV}) \) | \( T_h(\text{GeV}) \) | \( k \) |
|----------------------|----------------------------|-----------------|--------|
| 45.0                | 0.156                      | 0.709           | 3.23   |
| 36.2                | 0.154                      | 0.718           | 3.23   |
| 30.5                | 0.152                      | 0.705           | 3.18   |
| 23.2                | 0.150                      | 0.699           | 3.18   |
| 16.1                | 0.148                      | 0.678           | 3.15   |
| 9.8                 | 0.148                      | 0.655           | 3.12   |
| 4.3                 | 0.150                      | 0.613           | 3.11   |

**Table IV:** Characteristic temperatures and falloff \( k \) of the TMD of charged particles in \( XeXe \) collisions at \( \sqrt{s_{nn}} = 5.44 \) TeV for several centrality classes.

| \( dN_{ch}/d\eta \) | \( T_{th}(\text{GeV}) \) | \( T_h(\text{GeV}) \) | \( k \) |
|----------------------|----------------------------|-----------------|--------|
| 1167                | 0.138                      | 0.617           | 3.36   |
| 939                 | 0.136                      | 0.637           | 3.36   |
| 706                 | 0.135                      | 0.636           | 3.33   |
| 478                 | 0.132                      | 0.620           | 3.27   |
| 315                 | 0.129                      | 0.612           | 3.21   |
| 198                 | 0.126                      | 0.611           | 3.19   |
| 118                 | 0.123                      | 0.619           | 3.18   |
| 65                  | 0.119                      | 0.602           | 3.13   |
| 32                  | 0.114                      | 0.615           | 3.17   |

**Table V:** Characteristic temperatures and falloff \( k \) of the TMD of charged particles in \( PbPb \) collisions at \( \sqrt{s_{nn}} = 5.02 \) TeV for several centrality classes.

| \( dN_{ch}/d\eta \) | \( T_{th}(\text{GeV}) \) | \( T_h(\text{GeV}) \) | \( k \) |
|----------------------|----------------------------|-----------------|--------|
| 1942                | 0.140                      | 0.598           | 3.41   |
| 1585                | 0.139                      | 0.606           | 3.39   |
| 1180                | 0.138                      | 0.612           | 3.36   |
| 786                 | 0.135                      | 0.613           | 3.31   |
| 512                 | 0.132                      | 0.614           | 3.27   |
| 318                 | 0.129                      | 0.606           | 3.22   |
| 183                 | 0.125                      | 0.606           | 3.19   |
| 96                  | 0.120                      | 0.602           | 3.17   |
| 45                  | 0.116                      | 0.574           | 3.10   |
The above equation can be now cast into

$$F(p_{\perp}) = \int_0^{\infty} dT_h G(T_h) \frac{1}{(1 + p_{\perp}^2/\gamma)^k} = \int_0^{\infty} dT_h G(T_h) \int_0^x dx W_p(x) e^{x p_{\perp}^2},$$

(7)

The above equation can be now cast into

$$F(p_{\perp}) = \int_0^{\infty} dx W(x) e^{-x p_{\perp}^2},$$

(8)

where now $W(x)$ can be approximated by

$$W(x) = N \frac{x^{k-1}}{(1 + x/\bar{x})^{k'}},$$

(9)

where $N$ is a normalization constant $\bar{x}$ is the scale of $x$ and $k' = k + 1/2$ is required to obtain Gaussian fluctuations at high $x$. The above equation has the correct asymptotic behavior for both low and high $x$. In this way the transverse momentum distribution (8) becomes

$$F(p_{\perp}) = N \bar{x}^k \Gamma(k) U(k; 1/2, \bar{x} p_{\perp}^2).$$

(10)

$U$ is the confluent hypergeometric function. Its asymptotic limits are for $p_{\perp} \to 0$

$$F(p_{\perp}) = C \left(1 - \frac{2\Gamma(k + 1/2)}{\Gamma(k)} \bar{x}^{1/2} p_{\perp}\right),$$

(11)

and for $p_{\perp} \to \infty$

$$F(p_{\perp}) = \frac{C'}{(\bar{x} p_{\perp}^2)^x}.$$  

(12)

At high $p_{\perp}$, a power like behavior is obtained with power $2k$, and at low $p_{\perp}$ the thermal behavior

$$\exp(-p_{\perp}/T_{th})$$

with

$$T_{th} = \frac{\Gamma(k)}{2\Gamma(k + 1/2)} \frac{1}{\bar{x}^{1/2}}.$$  

(13)

As the fluctuations encoded in the function $W(x)$ depend only on the scale $\bar{x}$ and $k$, we can say that the effective thermal temperature depends only on the scale $\bar{x}$ and its fluctuations. In Figure 5, we show the fit using the whole soft and hard spectrum (10) to $pp$ data at different energies. A good description is obtained in all cases.

### III. MULTIPLICITY DISTRIBUTION SCALES

In what concerns to the multiplicity distribution, we note that a gamma distribution on the number of partons is also obtained for events which have at least one high $p_{\perp}$ particle due to a hard parton collision. In fact, if $P(n)$ is the probability of having $n$ partons in a given collision, the probability $P_c(n)$ of having $n$ partons with at least one hard is

$$P_c(n) = \frac{n}{\langle n \rangle} P(n).$$

(14)

This selection procedure can be repeatedly applied forming the chain

$$P(n) \rightarrow n \langle n \rangle P(n) \rightarrow n^2 \langle n^2 \rangle P(n) \rightarrow n^k \langle n^k \rangle P(n).$$

(15)

Similarly, we also notice that a gamma distributed multiplicity density convoluted with a Poisson process, produces a negative binomial distribution (NBD) for the multiplicity, broadly used to describe the experimental
data. Namely

\[
\frac{\Gamma(n + k_n)}{\Gamma(n + 1) \Gamma(k_n)} \frac{\gamma_n^{k_n}}{(1 + \gamma_n)^{k_n + n}} = \int_0^\infty dN e^{-N} N^n \frac{\gamma_n}{N} \tag{16}
\]

\[
x (\gamma_n N)^{k_n - 1} \exp(-\gamma_n N) = \int_0^\infty dN P(n, N) W_n(N),
\]

where as before

\[
\frac{1}{k_n} = \frac{\langle N^2 \rangle - \langle N \rangle^2}{\langle N \rangle^2}, \quad \gamma_n \equiv \frac{k_n}{\langle N \rangle}. \tag{17}
\]

Since the mean multiplicity and the \( p_\perp \) distribution are related as follows

\[
\langle n \rangle = \int d^2 p_\perp \int dx f(p_\perp, x) W_p(x), \tag{18}
\]

the following relation between the two gamma distributions can be written

\[
W_p(x) = \frac{\gamma_n}{\gamma} W_n(x), \tag{19}
\]

with \( k = k_n + 2 \). The convolution of the gamma distribution with a Poisson distribution gives rise to a negative binomial distribution for the multiplicity distribution, see formula (15). As far as the gamma distribution is obtained for at least one hard parton, the resulting multiplicity distribution describes the multiplicity distribution events with at least one hard parton. For the rest of events another distribution is required. There are several fits to the \( pp \) data at different colliding energies using two negative binomial distributions \([37–40]\). Each of these distributions has two parameters, \( k \) fixing the fluctuations and \( \langle n \rangle \) the mean multiplicity. The two mean multiplicities can be seen as the two multiplicity scales corresponding to the two transverse momentum scales \( T_h \) and \( T_p \), and we can look for a relation between these scales similar to the relation (4).

According to formula (19), the parameter \( k \) of the gamma distribution is two units larger than the one corresponding to the gamma distributions on the number of partons, thus we expect that instead of the equation (4) for the \( p_\perp \) distribution, we should have

\[
\frac{\langle N_h \rangle}{k - 1} = \langle N_h \rangle. \tag{20}
\]

This equation must be seen with caution because in the case of multiplicities longitudinal momentum fluctuations add to the \( p_\perp \) fluctuations. In order to avoid these contributions we look at the data on small rapidity range. In Table VI we show the results of a fit \([37]\) using two negative binomial distributions to the multiplicity distributions of \( pp \) collisions in the pseudo-rapidity range \(|\eta| \leq 0.5\) for different energies. The comparison of columns two and four shows a reasonable agreement.

| \( \sqrt{s_{nn}} \) (TeV) | \( \langle N_h \rangle \) | \( \langle N_h \rangle \) | \( \langle N_h \rangle / (k - 1) \) |
|-------------------------|----------|----------|-----------------|
| 0.9                     | 2.1±1.9  | 5±4      | 1.9±1.5         |
| 2.76                    | 2.5±1.0  | 7±2      | 3.1±0.9         |
| 7                       | 3.6±1.4  | 12±3     | 5.8±1.4         |

### IV. DISCUSSION

In the left hand side of equation (4), \( T_h \) and \( k \) are parameters related to hard collisions and thus described by perturbative QCD. On the other hand \( T_p \) has to do with non-perturbative QCD. Thus, to some extent, equation (4) links perturbative and non-perturbative physics. The factor \( k + 1 \) in equation (4) determines the fall-off of the probability of having an additional hard parton normalized by the hard scale \( T_h^2 \). Such probability is just the variation of the hard transverse momentum distribution.

Relations between these two regimes have been recently put forward in different quantum problems \([33]\) and, more recently, have been suggested as entropy constraints in an entangled nucleon, relating the final state multiplicity of the fragmenting nucleon with the parton distribution function probed by hard processes in \( pp \) and \( ep \) collisions \([7]\). Testing the implications of this entanglement, the H1 collaboration has measured very recently the charged particle distribution in DIS at HERA \([41]\). The hadron entropy found in data does not confirm, however, these entanglement predictions.

It is convenient to discuss if this disagreement is related to the way in which the entanglement entropy of the nucleon has to be obtained. There is not, to our knowledge, a known way of computing from first principles the distribution of weights in the entangled nucleon. Cascade models not including saturation and non-linear evolution \([18]\) may not be sensitive enough to correctly describe the initial entanglement entropy. Thus the observed entropy in the multiplicity distribution of hadrons may not be a in an one-to-one correspondence with the entropy of the cascade.

However, we notice that the observed multiplicity distributions in data are well described by NBDs. We can devise then a way of reconstructing the cascade process under these phenomenological considerations. We may assume the weights of the entangled state to be formed according to a Poisson process, with a given mean value \( \lambda \). Then each of these partons gives rise to a cascade with a geometric distribution \([18]\), with parameter \( p = (\langle n \rangle / (\langle n \rangle + k_n)) (k_n = 1$ for a geometric distribution). For the case of DIS at the energies explored by the H1 collaboration \([41]\), the fall-off parameter \( k_n \) is large enough to transform back the final NBD to the initial
Poisson distributed multiplicity. Following this observation, \(p\) has to be small and the final state multiplicity mirrors the initial weight distribution instead of being geometrically distributed. At larger energies, however, \(p\) must increase so that the part of the cascade becomes more important.

At the level of the Von Neumann entropy, this energy evolution corresponds to the passage from a Poisson distributed entropy to a geometric/Gamma distributed entropy

\[
\log \langle n \rangle^{1/2} \rightarrow \log \langle n \rangle
\]

the fall-off \(k\) of the NBD interpolating between these two limits \([20]\). At LHC energies for \(pp\) collisions we expect that the parameter \(p\) is large enough that the entropy is dominated by the geometric term and \(S \propto \log \langle n \rangle\), as was pointed out in \([7]\). At even larger energies due to saturation of partons we expect to recover a behavior \(S \propto \log \langle n \rangle^{1/2}\). Accordingly, the number of micro-states is not anymore \(n\), but saturates as \(\sqrt{n}\) when the colliding energy increases, following the expectations of the glasma picture of the CGC \([42]\) or the string percolation model \([43]\).

\[\text{V. CONCLUSIONS}\]

Summarizing up, the analysis of the transverse momentum distributions of \(pp\), \(p\bar{p}\), \(XeXe\) and \(PbPb\) collisions at different RHIC and LHC energies and centralities together with Higgs production decaying into \(\gamma\gamma\) and 4l suggest that a hard collision provides an ultraviolet scale that quenches the spectrum by means of fluctuations of the hard scale. A simple relation between the effective temperature and the hard scales is obtained which is approximately satisfied in the different cases in study. In this way, a non-perturbative scale \(T_{th}\) has been related to two perturbative quantities, \(T_h\) and \(k\). A gamma distribution is found, in agreement with phenomenological descriptions, for the distribution of the hard scale as well as the number of partons. The normalized fluctuations of both distributions are related and give rise to a relation between the multiplicities of the soft and hard spectrum. These findings are in line with the possibility that a hard parton collision works as an ultraviolet cutoff producing a quench of the rest of the entanglement partons of the initial wave function. Such entanglement may be at the origin of the apparent thermalization of the colliding hadrons.

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