1 Introduction

In high energy nuclear collisions\footnote{1} we are interested to create and study the deconfined quark-gluon phase, believed to consist in the limit of very high energy densities of a liquid of quarks and gluons interacting perturbatively. We are motivated by the desire to recreate in the laboratory conditions akin to those prevailing in the first moments of the Early Universe, and in the recognition that we can study the properties of the excited, ‘melted’, vacuum state of strong interactions.

Since in the collision of large nuclei the highly dense state is formed for a rather short time\footnote{2}, one of the major challenges has been to identify suitable physical observables of deconfinement. A number of possible experimental signatures of the formation and properties of quark-gluon plasma (QGP) have been studied. This report addresses our recent advances in evaluating the relevance of strange\footnote{3},\footnote{4} and charm quark flavor in this quest\footnote{4}. Other major probes include the phenomenon of $J/\Psi$ suppression\footnote{5}, photons and dileptons\footnote{6}.

Strangeness is a very interesting diagnostic tool of dense hadronic matter\footnote{3}:

1) particles containing strangeness are found more abundantly in relativistic nuclear collisions than it could be expected based on simple scaling of $p$–$p$ reactions;

2) all strange hadrons have to be made in inelastic reactions, while light $u$, $d$ quarks are also brought into the reaction by the colliding nuclei;

3) because there are many different strange particles, we have a very rich field of observables with which it is possible to explore diverse properties of the source;

4) theoretical calculations suggest that glue–glue collisions in the QGP provide a sufficiently fast and thus by far a unique mechanism leading to an explanation of strangeness enhancement.

We begin by recalling these mechanisms of strangeness production. We then obtain in section \ref{strangenessproduction} the magnitude of running coupling constant and quark masses. Using these results we evaluate the thermal relaxation times of strangeness and charm as function of temperature in section \ref{thermalsystem}. We discuss the relevance of the flavor observable of QGP and describe how these relaxation times allow to compute the hadronic particle yields in section \ref{hadronization}. We close with a few general remarks about the relevance of our results.

2 Strangeness Production in QGP

We will now show how to evaluate using two particle collision processes flavor production\footnote{6} in thermal QGP. Ultimately, we will employ running QCD parameters $\alpha_s(\mu)$ and $m_i(\mu)$, $i = s, c$. While theory and experiment constrain now sufficiently the coupling strength $\alpha_s$, considerable uncertainty still remains in particular in regard of strange quark mass scale, as well as systematic uncertainty related to applications of QCD to soft (less than 1 GeV) processes.

The generic angle averaged two particle cross
section for (heavy) flavor production processes \( g + g \to f + f \) and \( q + \bar{q} \to f + f \), are:

\[
\bar{\sigma}_{gg \to ff} = \frac{2\pi\alpha_s^2}{3s} \left[ \left( 1 + \frac{4m_f^2}{s} + \frac{m_f^2}{s^2} \right) \cdot \tanh^{-1} W(s) - \left( \frac{7}{8} + \frac{31m_f^2}{8s} \right) W(s) \right], \tag{1}
\]

\[
\bar{\sigma}_{q\bar{q} \to ff} = \frac{8\pi\alpha_s^2}{27s} \left( 1 + \frac{2m_f^2}{s} \right) W(s), \tag{2}
\]

where \( W(s) = \sqrt{1 - 4m_f^2/s} \), and both the QCD coupling constant \( \alpha_s \) and flavor quark mass \( m_f \) will be in this work the running QCD parameters. In this way a large number of even-\( \alpha_s \) diagrams contributing to flavor production is accounted for.

What remains unaccounted for is another class of processes in which at least one additional gluon is present. In particular processes allowing the production of an additional soft gluon in the final state remains unaccounted for today. Leading diagrams contain odd powers of \( \alpha_s \) and their generic cross section is in general infrared divergent, requiring a cut-off which for processes occurring in matter is provided by the interactions (dressing) with other particles present. The process in which a massive ‘gluon’, that is a quasiparticle with quantum numbers of a gluon, decays into a strange quark pair, is partially included in the resummation that we accomplish in the present work. At the present time we do not see a systematic way to incorporate any residue of this and other effects, originating in matter surrounding the microscopic processes, as work leading to understanding of renormalization group equations in matter (that is at finite temperature and/or chemical potential) is still in progress.

### 3 Running QCD Parameters

To determine the two QCD parameters required, we will use the renormalization group functions \( \beta \) and \( \gamma_m \):

\[
\mu \frac{\partial \alpha_s}{\partial \mu} = \beta(\alpha_s(\mu)), \quad \mu \frac{\partial m}{\partial \mu} = -m \gamma_m(\alpha_s(\mu)). \tag{3}
\]

For our present study we will use the perturbative power expansion in \( \alpha_s \):

\[
\beta_{\text{pert}} = \alpha_s^2 \left[ b_0 + b_1 \alpha_s + b_2 \alpha_s^2 + \ldots \right],
\]

\[
\gamma_m = \alpha_s \left[ c_0 + c_1 \alpha_s + c_2 \alpha_s^2 + \ldots \right]. \tag{4}
\]

For the SU(3)-gauge theory with \( n_f \) fermions the first two terms (two ‘loop’ order) are renormalization scheme independent, and we include in our calculations the three ‘loop’ term as well, which is renormalization scheme dependent, evaluated in the MS-scheme. We have:

\[
b_0 = \frac{1}{2\pi} \left( 11 - \frac{2}{3} n_f \right), \quad b_1 = \frac{1}{4\pi^2} \left( 51 - \frac{19}{3} n_f \right), \tag{5}
\]

\[
b_2 = \frac{1}{64\pi^3} \left( 2857 - \frac{5033}{9} n_f + \frac{325}{27} n_f^2 \right), \tag{6}
\]

\[
c_0 = \frac{2}{\pi}, \quad c_1 = \frac{1}{12\pi^2} \left( 101 - \frac{10}{3} n_f \right), \quad c_2 = \frac{1}{32\pi^3} \left( 1249 - \frac{2216}{27} + \frac{160}{3} \zeta(3) n_f - \frac{140}{81} n_f^2 \right). \tag{7}
\]

The number \( n_f \) of fermions that can be excited, depends on the energy scale \( \mu \). We have implemented this using the exact phase space form appropriate for the terms linear in \( n_f \):

\[
n_f(\mu) = 2 + \sum_{i=s,c,b,t} \sqrt{1 - \frac{4m_i^2}{\mu^2}} \cdot \left( 1 + \frac{2m_i^2}{\mu^2} \right) \Theta(\mu - 2m_i),
\]

with \( m_s = 0.16 \text{ GeV}, m_c = 1.5 \text{ GeV}, m_b = 4.8 \text{ GeV} \). We checked that there is very minimal impact of the running of the masses in Eq. (7) on the final result, and will therefore not introduce that ‘feed-back’ effect into our current discussion. The largest effect on our solutions comes from the bottom mass, since any error made at about 5 GeV is amplified most. However, we find that this results in a scarcely visible change even when the mass is changed by 10% and thus one can conclude that the exact values of the masses and the nature of flavor threshold is at present of minor importance in our study.

We show the result of numerical integration for \( \alpha_s \) in the top portion of Fig. 1. First equation in (3) is numerically integrated beginning with an initial value of \( \alpha_s(M_Z) \). We use in this report the August 1996 World averaged \( \alpha_s(M_Z) = 0.118 \) for which the estimated error is \pm 0.003. This value is sufficiently precise to eliminate most of the uncertainty that has befallen much of our earlier studies. In addition, the thin solid lines present results for \( \alpha_s(M_Z) = 0.115 \) till recently the preferred result in some analysis, especially those at

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running mass equation is linear in $m$, it is possible to determine the universal quark mass scale factor
\[ m_r = m(\mu)/m(\mu_0). \]  
(9)
Since $\alpha_s$ refers to the scale of $\mu_0 = M_Z$, it is a convenient reference point also for quark masses. As seen in the bottom portion of Fig. 1, the change in the quark mass factor is highly relevant, since it is driven by the rapidly changing $\alpha_s$ near to $\mu \approx 1$ GeV. For each of the different functional dependences $\alpha_s(\mu)$ we obtain a different function $m_r$. The significance of the running of the charmed quark mass cannot be stressed enough, especially for thermal charm production occurring in foreseeable future experiments well below threshold, which amplifies the importance of exact value of $m_c$.

Given these results, we find that for $\alpha_s = 0.118$ and $m_s(M_Z) = 90 \pm 18$ MeV a low energy strange quark mass $m_s(1$ GeV) $\approx 200 \pm 40$ MeV, in the middle of the standard range $100 < m_s(1$ GeV) $< 300$ MeV. Similarly we consider $m_c(M_Z) = 700 \pm 50$ MeV, for which value we find the low energy mass $m_c(1$ GeV) $\approx 1550 \pm 110$ MeV, at the upper (conservative for particle production yield) end of the standard range $1 < m_c(1$ GeV) $< 1.6$ MeV. There is another (nonperturbative) effect of mass running, related to the mass at threshold for pair production $m_{i\text{th}}$, $i = s, c$, arising from the solution of:
\[ m_{i\text{th}}/m_i(M_Z) = m_r(2m_{i\text{th}}). \]  
(10)
This effect stabilizes strangeness production cross section in the infrared: below $\sqrt{s} = 1$ GeV the strange quark mass increases rapidly and the threshold mass is considerably greater than $m_s(1$ GeV). We obtain the threshold values $2m_{s\text{th}} = 611$ MeV for $\alpha_s(M_Z) = 0.118$ and $2m_{c\text{th}} = 566$ MeV for $\alpha_s(M_Z) = 0.115$. Both values are indicated by the black dots in Fig. 1. For charm, the running mass effect plays differently: since the mass of charmed quarks is listed in tables for $\mu = 1$ GeV, but the value of the mass is above 1 GeV, the production threshold mass is smaller than expected (i.e., listed value). For $m_c(M_Z) = 700$ MeV the production threshold is found at $\sim 2m_{c\text{th}} \approx 2.3$ GeV rather than 3.1 GeV that would have been expected for the $m_c(1$ GeV). This reduction in threshold enhances thermal production of charm, especially so at low temperatures.
4 Strangeness and Charm Thermal Relaxation Times

The thermal average of the cross section is the invariant production rate per unit time and volume:

\[ A_s = A_{gg} + A_{u\bar{u}} + A_{d\bar{d}} + \ldots \]

\[ = \int_{4m_s^2}^{\infty} ds 2s \delta(s - (p_1 + p_2)^2) \int \frac{d^3p_1}{(2\pi)^3 E_1} \int \frac{d^3p_2}{2(2\pi)^3 E_2} \left[ \frac{1}{2} g_s^2 f_g(p_1) f_g(p_2) \bar{\sigma}_{gg}(s) + \ldots \right]. \tag{11} \]

The dots indicate that other mechanisms may contribute to strangeness production. The particle distributions \( f_i \) are in our case thermal Bose/Fermi functions (for fermions with \( \lambda_q = 1.5 \)), and \( g_1 = 6 \), \( g_8 = 16 \). For strangeness production \( n_t = 2 \), and for charm production \( n_t = 3 \).

\( \tau_s \) is from the invariant rate we obtain the strangeness relaxation time \( \tau_s \) shown in Fig. 2, as function of temperature:

\[ \tau_s = \frac{1}{2} \rho_s^{\infty}(\bar{m}_s). \tag{12} \]

Note that here unaccounted for processes, such as the above mentioned odd-order in \( \alpha_s \) would add to the production rate incoherently, since they can be distinguished by the presence of incoming/outgoing gluons. Thus the current calculation offers an upper limit on the actual relaxation time, which may still be smaller. In any case, the present result suffices to confirm that strangeness will be very near to chemical equilibrium in QGP formed in collisions of large nuclei.

We show in Fig. 2 also the impact of a 20% uncertainty in \( m_s(M_Z) \), indicated by the hatched areas. This uncertainty is today much larger compared to the uncertainty that arises from the recently improved precision of the strong coupling constant determination. We note that the calculations made at fixed values \( \alpha_s = 0.5 \) and \( m_s = 200 \text{ MeV} \) (dotted line in Fig. 2) are well within the band of values related to the uncertainty in the strange quark mass.

Since charm is somewhat more massive compared to strangeness, there is still less uncertainty arising in the extrapolation of the coupling constant. Also the systematic uncertainty related to the soft gluons (odd-\( \alpha_s \)) terms are smaller, and thus the relaxation times \( \tau_c \) we show in Fig. 2 are considerably better defined compared to \( \tau_s \). There is also less relative uncertainty in the value of charm mass. We also show in Fig. 2 (dotted lines) the fixed \( m_c \), \( \alpha_s \) results with parameters selected to border high and low \( T \) limits of the results presented. It is difficult to find a good comparative behavior of \( \tau_c \) using just one set of \( m_c \) and \( \alpha_s \). This may be attributed to the importance of the mass of the charmed quarks, considering that the threshold for charm production is well above the average thermal collision energy, which results...
in emphasis of the effect of running charm mass. In the high $T$-limit the choice (upper doted line in Fig. 3) $m_c = 1.5$ GeV, $\alpha_s = 0.4$ is appropriate, while to follow the result at small $T$ (lower doted line in Fig. 2) we take a much smaller mass $m_c = 1.1$ GeV, $\alpha_s = 0.35$.

We recall that the equilibrium distribution is result of Boltzmann equation description of two body collisions. Thus the mass arising in the equilibrium density $\rho_{\infty}$ in Eq. (12) is to be taken at the energy scale of the average two parton collision. We adopt for this purpose a fixed value $\tilde{m}_s = 200$ MeV, and observe that in the range of temperatures here considered the precise value of the mass is insignificant, since the quark density is primarily governed by the $T^3$ term in this limit, with finite mass correction being $O(10\%)$. The situation is less clear for charm relaxation, since the running of the mass should have a significant impact. Short of more complete kinetic treatment, we used $m_c \approx 1.5$ GeV in order to establish the reference density $\rho_{\infty}$ in Eq. (12).

5 QGP Flavor Observable

We will indicate in this section how the study of flavor production impacts our understanding and diagnosis of the deconfined QGP phase. We recall first that there are two generic flavor observable which we can study analyzing experimental data:

- **yield of strangeness/charm:** once produced in hot early QGP phase, strangeness/charm is not reannihilated in the evolution of the deconfined state towards freeze-out, and thus the flavor yield is characteristic of the initial, most extreme conditions;
- **phase space occupancy $\gamma_{s,c}$:** impacts distribution of flavor among final state particle abundances.

Given that the thermal equilibrium is established within a considerably shorter time scale than the (absolute) heavy flavor chemical equilibration, we can characterize the equilibration of the phase space occupancy by an average over the momentum distribution:

$$\gamma_i(t) = \frac{\int d^3p \int d^3x n_i(p, x, t)}{\int d^3p \int d^3x n_i^\infty(p, x)} , i = s, c. \quad (13)$$

The chemical equilibrium density is indicated by upper-script ‘$\infty$’. When several carriers of the flavor are present, as is the case in the confined phase, $n_i$ is understood to comprise a weighted sum.

In order to be able to compute the production and evoluntion of strangeness and charm flavor a more specific picture of the temporal evolution of dense matter is needed. Here, we will address specifically strangeness production in collisions at CERN-SPS, up to 200 A GeV per nucleon. We use a simple, qualitative description, simplified by the assumption that the properties of the hot, dense matter are constant across the entire volume (fireball model). We consider radial expansion to be the dominant factor for the evolution of the fireball properties such as temperature/energy density and lifetime of the QGP phase. Within the fireball model the expansion dynamics follows from two assumptions:

- the (radial) expansion is entropy conserving, thus the volume and temperature satisfy:

$$V \cdot T^3 = \text{Const.} \quad (14)$$

- the surface flow velocity is given by the sound velocity in a relativistic gas

$$v_t = 1/\sqrt{3}. \quad (15)$$

This leads to the explicit forms for the radius of the fireball and its average temperature:

$$R = R_{in} + \frac{1}{\sqrt{3}}(t - t_{in}), \quad T = \frac{T_{in}}{1 + (t - t_{in})/\sqrt{3} R_{in}}. \quad (16)$$

The initial conditions for Pb–Pb:

- $\lambda_q = 1.6, \quad t_{in} = 1$ fm/c, $T_{in} = 320$ MeV; with $R_{in} = 4.5$ fm for $\eta = 0.5; \quad R_{in} = 5.2$ fm for $\eta = 0.75$,

- and for S–Pb/W:

- $\lambda_q = 1.5, \quad t_{in} = 1$ fm/c, $T_{in} = 280$ MeV; with $R_{in} = 3.3$ fm for $\eta = 0.35; \quad R_{in} = 3.7$ fm for $\eta = 0.5$, have been determined such that the energy per baryon is given by energy and baryon flow, and the total baryon number is $\eta(A_1 + A_2)$, as stopped in the interaction region. The radius shown above are for zero impact parameter. For this, equations of state of the QGP are needed, and we have employed our model which the perturbative correction to the number of degrees of freedom were incorporated along with thermal particle masses.

In the fireball in every volume element we have:

$$n_i(p; t) = \gamma_i n_{c}^\infty(p; T, \mu_s). \quad (17)$$
In this limit and allowing for the detailed balance reactions, thus re-annihilation of flavor, the yield is obtained from the equation:

$$\frac{dN_s(t)}{dt} = V(t)A_s\left[1 - \gamma_s^2(t)\right]. \quad (17)$$

Allowing for dilution of the phase space density in expansion, we derive from Eq. (17) an equation describing the change in $\gamma_s(t)$:

$$\frac{d\gamma_s}{dt} = \left(\gamma_s \frac{Tm_s}{T^2} \frac{d}{dx} \ln x^2 K_2(x) + \frac{1}{2\gamma_s} [1 - \gamma_s^2]\right). \quad (18)$$

Here $K_2$ is a Bessel function and $x = m_s/T$. Note that even when $1 - \gamma_s^2 < 1$ we still can have a positive derivative of $\gamma_s$, since the first term on the right hand side of Eq. (18) is always positive, both $T$ and $d/dx(x^2k_2)$ being always negative. This shows that dilution due to expansion effects in principle can make the value of $\gamma_s$ rise above unity.

Given the relaxation constant $\tau_s(T(t))$, these equations can be integrated numerically, and we can obtain for the two currently explored experimental systems the values of the two observables, $\gamma_s$ and $N_s/B$, which are given in Table 1.

For S–Ag collisions at 200 A GeV a recent evaluation of the specific strangeness yield leads to $N_s/B|_{\text{exp}} = 0.86 \pm 0.14$ (see table 4 of Ref. 2). Our earlier analysis of the WA85 data yields $\gamma_s = 0.75 \pm 0.1$ for S–W interactions. Both these results are in good agreement with the theoretical result shown in the table, favoring the 50% stopping case for S–W. The analysis of the experimental Pb–Pb data is in progress.

As we can see in Table 1, there is a considerable uncertainty due to the unknown mass of strange quark. On the other hand, inspecting the yield of strange quarks per baryon $N_s/B$ there seem to be very little dependence on the stopping fractions. This insensitivity to the reaction mechanism coupled to a visible sensitivity to strange quark mass suggests that additional insight may be ultimately gained about the strange quark mass from the study of strangeness enhancement in relativistic heavy ion collisions. Thus we show in Fig. 4 as function of $m_s(M_Z)$ for the two different stopping fractions the resulting strangeness yield per baryon, $N_s/B$. The range of strange quark mass shown corresponds to the allowed range $45 < m_s(M_Z) < 135$ MeV ($100 \leq m_s(1 \text{ GeV}) \leq 300$ MeV). The known systematic and statistical error is indicated by the divergence of the different curves, and is particularly small for $N_s/B$.

Before we can use these results to obtain a more reliable estimate of strange quark mass, we will need to understand other sources of systematic uncertainty, e.g., those associated with unknown mechanisms of strangeness production and

Table 1: $\gamma_s$ and $N_s/B$ in S–W at 200 A GeV and Pb–Pb at 158 A GeV for different stopping values of baryonic number and energy $\eta_B = \eta_B$; computed for strange quark mass $m_s(1\text{ GeV}) = 200 \pm 40$ MeV, $\alpha_s(M_Z) = 0.118$.

| $E_{\text{lab}}$ | S–W at 200 A GeV | Pb–Pb at 158 A GeV |
|-----------------|------------------|---------------------|
| $\eta_B = \eta_B$ | 0.35             | 0.5                 |
| $\gamma_s$      | 0.55 ± 0.14      | 0.65 ± 0.15         |
| $N_s/B$         | 0.67 ± 0.16      | 0.70 ± 0.16         |

Figure 4: Phase space occupancy $\gamma_s$ and yield of strange quarks per baryon $N_s/B$ as function of strange quark mass $m_s(M_Z)$ for Pb–Pb collision system at 158 A GeV, for the two stopping fractions and $\alpha_s(M_Z) = 0.118 \pm 0.003$: thick solid lines, $\eta = 75\%$, with long dashed lines indicating the $\alpha_s(M_Z)$ uncertainty, thick dashed lines for $\eta = 50\%$, with dotted lines indicating the $\alpha_s(M_Z)$ uncertainty.
to improve the model of QGP we are employing. However, these results are so strongly dependent on $m_s$ and so little on other quantities, that we can be optimistic that one day strange and charm flavor mass may be obtained from the flavor yields seen in nuclear collisions.

Assuming that the model proposed has been effectively tested at 200 A GeV S–W/Pb collisions, we compute the strangeness yield and phase space occupancy as function of energy. The results are shown for the case of Pb–Pb collisions. We show these results in Fig. 5 as function of laboratory collision energy, stressing that if at low energies the QGP phase is not encountered, we would expect to see a drop in strangeness yield beyond the expectations here presented.

These results allow us to evaluate the strange (anti)baryon yields from QGP as function of collision energy. We note that at fixed $m_\perp$ the medium dependent factor controlling the abundance of hadrons emerging from the surface of the deconfined region is related to the chemical conditions in the source, and for strange quarks, there is also the occupancy factor $\gamma_s$ to be considered:

$$n_h|_{m_\perp} = e^{-m_\perp/T} \prod_{k \in h} \gamma_k \lambda_k .$$  \hspace{1cm} (19)

The strange quark fugacity is in deconfined phase unity, while the light quark fugacity evolution with energy of colliding ions follows from our earlier studies \cite{1}. In Fig. 6 we have normalized all yields at $E_{\text{Lab}} = 158$ A GeV. Remarkably, all antibaryon yields (left hand side of Fig. 6) cluster together (solid lines: (anti)nucleons, long dashed: (anti)hyperons, short-dashed: (anti)cascades, and dotted: (anti)omegas), thus as long as the QGP phase is formed, ratios of rare multistrange antibaryons should not change significantly while the collision energy is reduced, until the QGP formation is disrupted. It should be noted that the yield of $\Omega$ remains appreciable, all the way even at very small energies — this is the case as long as these particles are produced by the deconfined phase, rather than in individual hadronic interactions. For baryons (right hand side of Fig. 6) there is considerable differentiation of the yield behavior: the reference yield of nucleons rises slightly to compensate for the drop in the yield of other strange baryons, which decreases substantially with energy.

6 Final Remarks

In conclusion, we can once more remind ourselves about the two generic strangeness observables. The relative total abundance of strangeness is most related to the initial condition, the ‘hotter’ the initial state is, the greater the production rate, and thus the final state relative yield, to be measured with respect to baryon number or global particle multiplicity (entropy). The phase space occupancy of strangeness $\gamma_s$ depends aside of the initial production rate, on the final state dilution characterized by dynamics of the expansion and the freeze-out temperature. Our results show that we can use certain features of strange particle production to see the formation of the deconfined state and to study some QCD properties.
and parameters. Our results suggest that already at present energies deconfinement is attained, and
we have explored a number of features as function of collision energy in order to see if a more syste-
matic study is capable to confirm this conclusion.

Using QCD renormalization group methods we have studied the s and c flavor chemical equi-
librium relaxation times. We have shown that the newly measured QCD coupling constant comprises
sufficiently small uncertainty to allow precise eval-
uation of strangeness production at and below 1
GeV energy scale. Our study has further proven
that it is essential to incorporate in the evaluation
of flavor production rates both running coupling
constant and running mass.

We find that running of the QCD parameters
is of major significance, since, e.g., the effective
charm production mass is considerably reduced,
seen on the scale of available thermal energies. We
found considerable enhancement of charm produc-
tion for temperatures applicable at SPS collision
energy, compared to fixed mass results. While
classical experiences at low temperature \( T \approx 200
\text{MeV} \) a 100 times slower approach to chemical
equilibrium compared to strangeness, for temperatures
of about 500 MeV, as may apply to the conditions
generated at LHC or perhaps even RHIC collider,
\( \tau_c \to 30 \text{fm} \), which is within factor two of the ex-
pected maximum lifespan of the deconfined state.
Thus our calculations suggest that there will be a
significant abundance of thermal charm in nuclear
interactions at \( T \approx 500 \text{MeV} \), while charm
may play a similar role in the diagnosis of the
‘hot’ \( T \approx 250 \text{MeV} \) case, and charm equilibrium appears within reach
of the extreme conditions possibly arising at LHC.

Our here presented results imply that in key
features the strange particle production results ob-
tained at 160–200 A GeV, are consistent with the
QGP formation hypothesis. However, in order
to ascertain the possibility that indeed the QGP
phase is already formed today, a more systematic
experimental exploration as function of collision
energy of the different observable is required,
for which purpose we also have explored the collision
energy dependence of the most characteristic strange particle features expected from the QGP
phase.

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