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Thermal Strangeness and Charm in QGP

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Abstract. Using QCD methods we re-evaluate strangeness and charm production in a thermal QGP fireball. QCD renormalization group is employed to evaluate running of the parameters $\alpha_s(\mu)$ and $m_s(\mu)$. We resum even- α_s Feynman-diagrams involving two particles in initial and final states. We obtain thermal relaxation times in a simple dynamical description of the fireball and we use these results to study two generic strangeness observables as function of collision energy: specific (per baryon) strangeness yield $\langle \bar{s} \rangle / B$, and phase space occupancy γ_s . From these two quantities, knowing the hadronization process in rough detail, we obtain the final state strange particle abundances.

1. Introduction

The thermal production of heavy flavor is a considerably slower process compared to the multitude of different reactions possible in relativistic nuclear collisions. However, the relaxation of strangeness population occurs in the deconfined phase within a time span comparable to the life time of the high density space-time region formed in such collisions [1, 2]. Consequently, quantitative understanding of QGP based strangeness production is a key element in our quest to use strange particle yields for diagnosis of the formation and the exploration of the properties of the deconfined phase. We report here briefly on the progress we have made [3, 4] in the past year in improving the understanding of the main features of strangeness production. We also apply our methods to the problem of thermal charm production.

Our objective has been to eliminate the need to use of an arbitrary coupling constant in the two particle collision cross section of strangeness and charm production, noting that the rate of strangeness production is proportional to α_s^2 and that the values between $0.2 < \alpha_s < 0.6$ were used in literature. This means that the different groups were using an order of magnitude different rates, no wonder that sometimes opposing views have arisen as to the magnitude and relevance of some results.

While the first order free space (perturbative vacuum) strangeness production processes at fixed values of $\alpha_s = 0.6$ and $m_s = 160-220$ MeV, have been considered

for some time [1], certain nonperturbative effects were more recently explored. The thermal production rates in medium, incorporating temperature dependent non-perturbative particle masses [5] have lead to the total strangeness production rate which was found little changed compared to the free space rate. This finding was challenged [6], but a more recent reevaluation of this work [7] confirmed that the rates obtained with perturbative glue-fusion processes are describing adequately the strangeness production rates in QGP, with a strangeness relaxation constant $\tau_s \simeq 3 \text{ fm}/c$ for the here relevant temperature range $T \simeq 250 \text{ MeV}$. For further discussion of the situation prior to Summer 1995 we refer to reference [8].

Once strangeness production rates in QGP are understood, an expansion model of the QGP phase takes us to the hadronization phase transition. Here some very crucial and little understood processes take place determining the primary hadronic particle abundances. Irrespective of the subsequent evolution, small rates in HG for strangeness annihilation assure that total strangeness abundance remains unchanged. Thus in principle one can observe in the final state, the ‘deepest’ history of strangeness production. However, a measurement of total strangeness yield is very difficult, and it has been reported only for the case of 200A GeV S–Ag collisions [9].

Individual strange particles become direct probes of the QGP phase if their production should occur prior to the global hadronization and/or they can escape without reequilibration with the global particle abundance. Such a reaction picture was proposed for high $m_\perp$ strange (anti)baryons [10], and it is possible that it applies to all hadrons, if the QGP phase is explosively disintegrating [11]. Because in the deconfined phase the strange and antistrange quarks occupy identical phase space volumes, the strange quark chemical fugacity $\lambda_s = 1$. In the HG at finite baryon density the equivalent chemical fugacity can take this value only in an exceptional condition. In view of the fact that strange (anti)baryon abundances observed display $\lambda_s = 1$ one is lead to argue that they indeed emerge directly from the deconfined phase, and thus reflect in their properties on the situation at the breakup of this state. This allows to describe the energy dependence of the strange (anti)baryon production employing the QGP chemical properties derived from the dynamical properties of the heavy ion collision [12] and the strange phase space evolution.

2. Flavor Production

The generic angle averaged cross sections for (heavy) flavor s , $\bar{s}$ production processes $g + g \rightarrow s + \bar{s}$ and $q + \bar{q} \rightarrow s + \bar{s}$, are [13]:

$$\bar{\sigma}_{gg \rightarrow s\bar{s}}(s) = \frac{2\pi\alpha_s^2}{3s} \left[\left(1 + \frac{4m_s^2}{s} + \frac{m_s^4}{s^2} \right) \tanh^{-1} W(s) - \left(\frac{7}{8} + \frac{31m_s^2}{8s} \right) W(s) \right], \quad (1)$$

$$\bar{\sigma}_{q\bar{q} \rightarrow s\bar{s}}(s) = \frac{8\pi\alpha_s^2}{27s} \left(1 + \frac{2m_s^2}{s} \right) W(s), \quad (2)$$

where $W(s) = \sqrt{1 - 4m_s^2/s}$, and both the QCD coupling constant α_s and strange quark mass m_s will be in this work the running QCD parameters taken at the energy scale $\mu \equiv \sqrt{s}$. In this way a large number of even- α_s diagrams contributing to the cross section is accounted for.

We employ the state of the art knowledge of the QCD renormalization group functions:

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

These functions β and γ_m are today known for the SU(3)-gauge theory with n_f fermions, but only in a perturbative power expansion in α_s ; three leading terms are commonly given for β and two for the γ_m [14]:

$$\beta^{\text{pert}} = \alpha_s^2 [b_0 + b_1 \alpha_s + b_2 \alpha_s^2 + \dots], \quad \gamma_m^{\text{pert}} = \alpha_s [c_0 + c_1 \alpha_s + \dots], \quad (4)$$

with

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

$$c_0 = \frac{2}{\pi}, \quad c_1 = \frac{1}{12\pi^2} \left(101 - \frac{10}{3} n_f \right). \quad (6)$$

The b_2 -coefficient in Eq. (4) is renormalization scheme dependent; it is known, e.g., in the modified minimum subtraction dimensional renormalization $\overline{MS}$ scheme:

$$b_2^{\overline{MS}} = \frac{1}{64\pi^3} \left(2857 - \frac{5033}{9} n_f + \frac{325}{27} n_f^2 \right). \quad (7)$$

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

$$n_f(\mu) = 2 + \sum_{i=s,c,b,t} \sqrt{1 - \frac{4m_i^2}{\mu^2}} \left(1 + \frac{2m_i^2}{\mu} \right) \Theta(\mu - 2m_i), \quad (8)$$

with $m_s = 0.16$ GeV, $m_c = 1.5$ GeV, $m_b = 4.8$ GeV. We checked that there is very minimal impact of the running of the masses in Eq. (8) 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. Thus we will present also (the scarcely visible) result obtained with $m_b = 4.4$ GeV, a sometimes preferred value.

Since we consider values $\alpha_s \rightarrow 1$ and recent work suggests that Pade approximant in QCD expression could improve the precision and enlarge the circle of convergence of the perturbative expansion [15] we have used the Pade-approximant for the β -function in form:

$$\beta \rightarrow \beta^{(0,2)} \equiv \alpha_s^2 b_0 \frac{1}{1 - u\alpha_s + v\alpha_s^2} = \alpha_s^2 b_0 \frac{b_0}{b_0 - b_1 \alpha_s + (b_1^2/b_0 - b_2) \alpha_s^2}. \quad (9)$$

Integration of the renormalization equations is facilitated by this modification since it turns out that the numerical solutions are stabilized in that way.

Fig.1. $\alpha_s(\mu)$ (top portion) and $m_r(\mu) = m(\mu)/m(M_{Z^0})$ (bottom portion) as function of energy scale μ . Thick lines correspond to initial value $\alpha_s(M_{Z^0}) = 0.105$, short-dashed line to the central standard case $\alpha_s(M_{Z^0}) = 0.115$, thin line is for the initial value $\alpha_s(M_{Z^0}) = 0.12$. Long-dashed line was obtained with the bottom flavor threshold reduced from 4.8 GeV to 4.4 GeV for $\alpha_s(M_{Z^0}) = 0.12$. In lower section the dots indicate the strangeness pair production thresholds for $m_s(M_{Z^0}) = 90$ MeV, while crosses indicate charm pair production thresholds for $m_c(M_{Z^0}) = 700$ MeV.

First equation in (3) is numerically integrated beginning with initial value of $\alpha_s(M_Z)$. We considered initially the values $\alpha_s(M_Z) = 0.12; 0.115; 0.105$. The middle value is the current best accepted value [14], other two limit the range possible, considering in particular fits to the line-shape of Z^0 when certain anomalous decays are not considered [16]. The smallest value leads for us to a lower limit on

the production rate of strangeness in QGP. We show the result in Fig. 1; thick lines correspond to the lowest limit, short-dashed line to the central standard case, the thin line is a reasonable upper limit on α_s . The long dashed lines just visible for large values of α_s show the impact of the changed mass for the bottom threshold, as mentioned above. Our central result corresponds exactly to the one considered in Ref. [14].

With $\alpha_s(\mu)$ from the solutions described above we integrate the running of the quark masses, the second equation in (3). Because the running mass equation is linear in m , it is possible to determine the universal multiplicative quark mass scale factor

$$m_r = m(\mu)/m(\mu_0). \quad (10)$$

Since α_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 α_s near to $\mu \simeq 1$ GeV. For each of the two different functional dependences $\alpha_s(\mu)$ we obtain a different function m_r .

For $m_s(M_{Z^0}) = 90$ MeV we find $m_s(1 \text{ GeV}) \simeq 160$ and 215 MeV for our two choices of $\alpha_s(M_{Z^0})$. We will investigate the impact of 20% uncertainty in the mass of the strange quark. Thus we explore the range $130 \leq m_s(1 \text{ GeV}) \leq 260$ MeV for the strange quark mass, within the standard range $100 < m_s(1 \text{ GeV}) < 300$ MeV.

The strangeness production cross section is subject to an implicit infrared stabilization: below $\sqrt{s} = 1$ GeV the strange quark mass increases rapidly and the threshold mass m_s^{th} for the pair production defined by the solution of:

$$m_s^{\text{th}}/m_s(M_{Z^0}) = m_r(2m_s^{\text{th}}) \quad (11)$$

is considerably greater than $m_s(1 \text{ GeV})$. We obtain the threshold values $2m_s^{\text{th}} = 470$ MeV for the smaller α_s option (thick lines) and $2m_s^{\text{th}} = 740$ MeV for the higher option (thin line). Both values are indicated by the black dots in Fig. 1.

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

$$\begin{aligned} A_s &\equiv A_{gg} + A_{u\bar{u}} + A_{d\bar{d}} = \sum_{AB} \langle \sigma v_{AB} \rangle_T \rho_A^\infty \rho_B^\infty \\ &= \int_{4m_s^2}^{\infty} ds 2s \delta(s - (p_A + p_B)^2) \int \frac{d^3 p_A}{(2\pi)^3 2E_A} \int \frac{d^3 p_B}{(2\pi)^3 2E_B} \\ &\quad \times \left[\frac{1}{2} g_g^2 f_g(p_A) f_g(p_B) \overline{\sigma}_{gg}(s) + n_f g_q^2 f_q(p_A) f_q(p_B) \overline{\sigma}_{q\bar{q}}(s) \right]. \end{aligned} \quad (12)$$

where the particle distributions f_i are in our case thermal Bose/Fermi functions, for fermions with $\lambda_q = 1.5$, and $n_f = 2$, $g_q = 6$, $g_g = 16$. From the invariant rate we obtain the strangeness relaxation time shown in Fig. 2 as function of temperature:

$$\tau_s \equiv \frac{1}{2} \frac{\rho_s^\infty}{(A_{gg} + A_{qq} + \dots)}, \quad (13)$$

where the dots indicate that other mechanisms may contribute to strangeness production, reducing the relaxation time obtained considering the processes of gluon and quark fusion.

Fig.2. QGP strangeness relaxation time. Upper limit for $\alpha_s(M_{Z^0}) = 0.105$: thick solid line; lower limit for $\alpha_s(M_{Z^0}) = 0.12$: thin solid line, note dashed trailing line for reduced m_b threshold. Dotted: comparison results for fixed $\alpha_s = 0.6$ and $m_s = 200$ MeV. Hatched areas — effect of variation of strange quark mass $m_s(M_{Z^0}) = 90$ MeV by $\pm 20\%$.

Since the solid line is the lower limit on the cross section, here it is the upper limit on the relaxation time. We recall that the equilibrium distribution is result of Boltzmann equation description of two body collisions. Thus the mass arising in the equilibrium density ρ_s^∞ in Eq. (13) is to be taken at the energy scale of the average two parton collision. We adopt for this purpose a fixed value $m_s = 200$ 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 $\mathcal{O}(10\%)$.

We show in Fig. 2 also the impact of a 20% uncertainty in the strange quark mass, indicated by the hatched areas. This uncertainty is of similar magnitude as the uncertainty that arises from the current knowledge of the strong coupling constant. We note that the calculations made [1] at fixed values $\alpha_s = 0.6$ and $m_s = 200$ MeV (dotted line in Fig. 2) agree well with current study, which provides a measure of allowed ranges for these parameters. In particular it is not possible

to use a fixed value of $\alpha_s < 0.4$ for the study of strangeness production.

To obtain the relaxation time for charm production we first note that the running mass effect also is important here. Remarkably, since the mass of charmed quarks is listed in tables for $\mu = 1$ GeV ($m_c(1 \text{ GeV}) = 1.5$ GeV), and the value of the mass is above 1 GeV, the production threshold mass is effectively smaller than the value listed. We take the value at $m_c(\mu = M_Z) = 700$ MeV, which choice assures given $\alpha_s(\mu)$ that $m_c(1 \text{ GeV}) \simeq 1.5$ GeV. The drop in m_c at the production threshold, $2m_c^{\text{th}} \simeq 2.5$ GeV rather than 3 GeV has important consequences for the rate of thermal charm production, since it is so difficult to find, at temperatures of interest, collisions at these high $\sqrt{s}$.

Fig.3. Solid lines: thermal charm relaxation constant in QGP, calculated for running $\alpha_s(M_{Z^0}) = 0.12; 0.105$, $m_c(M_Z) = 700$ MeV. Dotted line: for fixed $m_c = 1.1$ GeV, $\alpha_s = 0.35$; short-long dashed line: for fixed $m_c = 1.5$ GeV, $\alpha_s = 0.4$.

In Fig. 3 we show the charm relaxation constant, see Eq. (13), for the established α_s values — because of the different production threshold of the charm for the different α_s cases, the curves for the two $\alpha_s(M_{Z^0}) = 0.12; 0.105$ considered, which lead to very similar result, do cross. We used $m_c \simeq 1.5$ GeV in order to establish the reference density ρ_c^∞ in Eq. (13). We also show in Fig. 3 the fixed m_c, α_s results with parameters selected to border high and low T limits of the results presented. We see that in the high T -limit the choice (short-long dashed line in Fig. 3) $m_c = 1.5$ GeV, $\alpha_s = 0.4$ is appropriate, while to follow the result at small T (dotted line in Fig. 3) we must take a much smaller mass $m_c = 1.1$ GeV, $\alpha_s = 0.35$.

3. QGP observables

We now can proceed to explore two generic (strangeness) observables as function of the impact parameter (baryon content) and collision energy:

- Specific (with respect to baryon number B) strangeness yield N_s/B
Once produced strangeness escapes, bound in diverse hadrons, from the evolving fireball and hence the total abundance observed is characteristic for the initial extreme conditions;
- Phase space occupancy γ_s
Strangeness freeze-out conditions at particle hadronization time t_f , given the initially produced abundance, determine the final state observable phase space occupancy of strangeness $\gamma_s(t_f)$.

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

$$\gamma_{s,c}(t) \equiv \frac{\int d^3p d^3x n_{s,c}(\vec{p}, \vec{x}; t)}{\int d^3p d^3x n_{s,c}^\infty(\vec{p}, \vec{x})}, \quad (14)$$

where $n_{s,c}$ is the sum over all heavy flavor containing particle densities, and should multi-strange/charmed objects be present, this sum contains the associated weight. $n_{s,c}^\infty$ is the same, but for the equilibrium particle densities. In QGP deconfined state, of course we have practically just the free quarks.

To compute the two generic observables introduced, a more specific picture of the temporal evolution is needed. We will presume, in a qualitative approach that

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

$$V \cdot T^3 = \text{Const.} \quad (15)$$

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

$$v_f = 1/\sqrt{3}. \quad (16)$$

This leads to the explicit forms:

$$R = R_{\text{in}} + \frac{1}{\sqrt{3}}(t - t_{\text{in}}), \quad T = \frac{T_{\text{in}}}{1 + (t - t_{\text{in}})/\sqrt{3}R_{\text{in}}}. \quad (17)$$

The initial conditions:

$$\begin{aligned} T_{\text{in}} &= 320 \text{ MeV}; \quad R_{\text{in}} = 5.6 \text{ fm}; \quad t_{\text{in}} = 1 \text{ fm}/c; \quad \lambda_q = 1.6; \quad \text{for Pb-Pb}, \\ T_{\text{in}} &= 280 \text{ MeV}; \quad R_{\text{in}} = 4.7 \text{ fm}; \quad t_{\text{in}} = 1 \text{ fm}/c; \quad \lambda_q = 1.5; \quad \text{for S-Pb/W}, \end{aligned}$$

are derived from our exploration of the equations of state of QGP [12]: the radius R_{in} has been determined such that the baryon number content in the fireball is 380 (Pb–Pb case) and 120 (S–Pb/W case) respectively, corresponding to zero impact parameter collisions. The energy/baryon content at given projectile energy is computed assuming that the same stopping governs energy and baryon number.

When we assume that the fireball is homogeneous in T and λ_q we can write:

$$n_s(\vec{p}; t) = \gamma_s n_s^\infty(\vec{p}; T, \mu_s). \quad (18)$$

In this limit and allowing for the back reactions, thus re-annihilation of strangeness, the strangeness yield is simply obtained solving the equation

$$\frac{dN_s(t)}{dt} = V(t) A_s [1 - \gamma_s^2(t)]. \quad (19)$$

While the equations describing the change in $\gamma_s(t)$ is:

$$\frac{d\gamma_s}{dt} = \left(\gamma_s \frac{\dot{T} m_s}{T^2} \frac{d}{dx} \ln x^2 K_2(x) + \frac{1}{2\tau_s} [1 - \gamma_s^2] \right). \quad (20)$$

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 γ_s , since the first term on the right hand side of Eq. (20) is always positive, both $\dot{T}$ and $d/dx(x^2 K_2)$ being always negative. This dilution due to expansion effects in principle can make the value of γ_s rise above unity.

We solve these equations numerically and show the result in Fig. 4 as function of fireball energy content converted to beam energy, assuming that the stopping of baryon number and energy is equal. To the left in Fig. 4 we present $\gamma_s(t_f)$ which depends on both the initial production and the enrichment of the strange phase space occupancy by the dilution effect, and thus somewhat on freeze-out conditions, here assumed to occur at $T_f = 140$ MeV. High values of γ_s should accompany low freeze-out temperature, provided that there has been extreme initial conditions allowing to produce strangeness. To the right we present the specific strangeness yield. We recall that for S–Ag collisions at 200A GeV a recent evaluation of the strangeness yield leads to $N_s/B = 0.86 \pm 0.14$ (see table 4 of Ref. [9]) which fits also nicely into the range of values shown at 200 GeV in Fig. 4.

Knowledge about strangeness evolution described in Fig. 4 allows us to compute 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 particles 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 γ_s to be considered:

$$n_h|_{m_\perp} = e^{-m_\perp/T} \prod_{k \in h} \gamma_k \lambda_k. \quad (21)$$

Fig.4. Strangeness occupancy $\gamma_s(t_f)$ (left) and specific strangeness yield per baryon N_s/B (right) as function of beam energy for central S-W/Pb collisions (dashed lines) and Pb-Pb collisions (solid lines) assuming $m_s(M_Z) = 90$ MeV, for $\alpha_s(M_{Z^0}) = 0.12$; 0.105 and three dimensional expansion of the fireball — the vertical bar corresponds to the value of γ_s found in S-W data analysis[11].

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 [12]. In Fig. 5. we have normalized all yields at $E_{\text{Lab}} = 158A$ GeV. Remarkably all antibaryons (left hand side of Fig. 5) seem to 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 multi strange antibaryons should not decrease while the energy is reduced, till the QGP formation is disrupted. It should be noted that the yield of $\bar{\Omega}$ 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. 5) there is considerable differentiation of the yield behavior: the reference yield of nucleons remains constant, as it is intuitively expected, but the disappearance of more strange baryons is faster.

4. Final Remarks

Using renormalization group methods and best current knowledge about the QCD properties, we have studied the flavour (s, c) production cross sections. Our results are in qualitative agreement with earlier work that employed $\alpha_s \simeq 0.6$ but considerable detailed differences arise, which justify the effort made here. We have also evaluated the thermal charm production and find that running of the QCD

parameters is of crucial significance, since the effective production mass is considerably reduced, seen on the scale of available thermal energies. Our study has further proven that it is essential to incorporate in the evaluation of flavor production rates both running mass *and* running coupling constant. It is noteworthy, in view of results shown in Fig. 3, that charm equilibrium appears within reach of the extreme conditions possibly arising at LHC.

Fig.5. Yield of antibaryons (left) and baryons (right) as function of heavy ion collision energy E_{Lab} . Abundance set to a common value (unity) at $E_{\text{Lab}} = 158A$ GeV. Solid lines nucleons, long dashed $\Lambda + \Sigma^0$, short dashed Ξ , dotted Ω .

We have explored the total strange particle yields as function of collision energy and impact parameter and found that the yield of strangeness per baryon is primarily determined by the initial thermal properties of the QGP fireball, and in particular by the initial temperature. In a wide energy range the specific strangeness yield rises linearly with the fireball CM-energy content, reaching close to one strange quark pair per baryon at about $E/B = 8-9$ GeV.

Our here presented and other results [4] imply that in key features the strange particle production results obtained at 160–200A 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 observables 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. We determined that there is considerable sensitivity of strangeness observables to both the QCD properties and dynamical properties of the QGP fireball.

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