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STRANGENESS PRODUCTION IN THE QUARK GLUON PLASMA*

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ABSTRACT

It is shown that perturbative QCD predicts abundant strange quark production in the plasma created in high energy nuclear collisions. Considering further the strange particle production in the hadronic gas phase, I show that the strangeness abundance in the plasma is 10-50 times higher as compared with the gas phase in similar thermodynamic conditions. Possible experiments leading to the identification of the plasma phase are described.

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1. INTRODUCTION

In nuclear collisions at high energies, qualitatively new physical phenomena may be studied. Among them the possible formation of quark-gluon plasma is the most outstanding novel physical phenomenon. The subject of this paper is the proposal[1] that this new form of matter may very likely be uniquely characterised by abundances and correlations among strange hadrons produced.

The basic question to answer by such observations concerns the intrinsic structure of the high energy density space region created in nuclear collision. It can consist either of individual hadrons (hadronic gas phase) or instead of quarks and gluons, deconfined over the region of space where the plasma state is formed[2]. Can strange particles tell the difference between these two à priori so different phases of hadronic matter, and in particular, can they tell us that the collision has lead to a momentary formation of the plasma state?

I assume that the reader of this report is largely familiar with the fundamental elements of statistical physics and will use these powerful methods throughout this report when discussing regions of space in which there has been substantial energy (density), of the order of 1 GeV/fm^3 or more, deposited in the course of the nuclear collisions. We recall, in particular, that the unhandy extensive variables, *viz.*, energy, baryon number, etc., are replaced by intensive quantities. To wit, the temperature T is a measure of energy per degree of freedom; the baryon chemical potential μ controls the mean baryon density.

The theoretical techniques required for the description of the different phases *viz.*, the hadronic gas and the quark-gluon plasma, allow for the formation of numerous hadronic resonances on the one side[3] (statistical bootstrap model) which then at sufficiently high energy density dissolve into the state consisting of their constituents. At this point we must appreciate the importance and help by a finite, i.e., non-zero temperature in reaching the transition to the quark-gluon plasma: to obtain a high particle density, instead of only compressing the matter (which as it turns out is quite difficult), we also heat it up; many pions are generated in a collision, allowing the transition to occur at moderate, even vanishing baryon density[4]. As qualitative discussion shows, and detailed investigations confirm[5], we are led to the conjecture of a first order phase transition in such an approach. This conjecture of Ref.1a been criticized, and only more recent lattice gauge theory calculations have led to more widespread acceptance of this phenomenon provided an internal SU(3) (colour) symmetry is used - SU(2) internal symmetry leads perhaps to a second order phase transition[6]. However, I would like to emphasize that the question whether the transition hadronic gas ↔ quark-gluon plasma is a phase transition (i.e., discontinuous) or continuous phase transformation will probably only be answered in actual experimental work, as all theoretical approaches suffer from approximations unknown in their effect. For example, in lattice gauge computer calculations, we establish the properties of the *lattice* and not of the continuous space in which we live.

Further development of this new field of research depends on the ability to observe plasma's creation and its detailed physical properties. It is quite difficult to insert a thermometer and to measure baryon density at $T = 150$ MeV and threefold or even higher nuclear compressions. We must either use only electromagnetically interacting particles[7] (photons, lepton pairs) in order to get them out of the plasma or study the heavy quark flavour abundance, in particular strangeness, generated in the collision[1]. To obtain a better impression of what is meant, imagine that strange quarks are very abundant in the plasma (and indeed they are!). Then, e.g. since (sss)-state is bound and stable in the hot perturbative QCD-vacuum, it could be the most abundant baryon to emerge from the plasma. I doubt that such an Omegasation of nuclear matter could leave any doubts about the occurrence of the phase transition. But even observation of an enhancement of the more accessible abundance of $\overline{\Lambda}$, may already be sufficient for our purposes.

I now explain in more detail why strange particle abundance are so helpful[1] in order to observe properties of quark-gluon plasma. First we note that, at a given temperature, the quark-gluon plasma will contain an equal number of strange (s) quarks and antistrange ($\bar{s}$) quarks, naturally assuming that the hadronic collision time is much too short to allow for light flavour weak interaction conversion to strangeness. Thus, assuming absolute chemical equilibrium in the quark plasma (see Section 3), we find the density of the strange quarks to be (two spins and three colours):

$$s/v \;=\; \bar{s}/v \;=\; 6\int \frac{d^3p}{(2\pi)^3}\,\frac{1}{e^{\sqrt{p^2+M^2}/T}+1} \approx 3\,\frac{TM^2}{\pi^2}\,K_2(M/T) \qquad (1.1)$$

(neglecting, for the time being, the perturbative corrections). We recall that the mass of the strange quarks, M, in the perturbative vacuum is believed to be of the order of 150 MeV. In Eq.(1.1) we were able to use the Boltzmann limit, when the phase space density of strangeness is not too high. Similarly, there is a certain light antiquark density ($\bar{q}$ stands for either $\bar{u}$ or $\bar{d}$):

$$q/V \;\hat{=}\; 6\int \frac{d^3p}{(2\pi)^3}\,\frac{1}{e^{|p|/T+\mu_q/T}+1} \approx e^{-\mu_q/T}\,T^3\,\frac{6}{\pi^2}\;, \qquad (1.2)$$

where the quark chemical potential is $\mu_q = \mu/3$ and μ is baryonic chemical potential. This exponent supresses the $q\bar{q}$ pair production. It reflects the chemical equilibrium between q-$\bar{q}$ and the presence of a light quark density associated with the net baryon number.

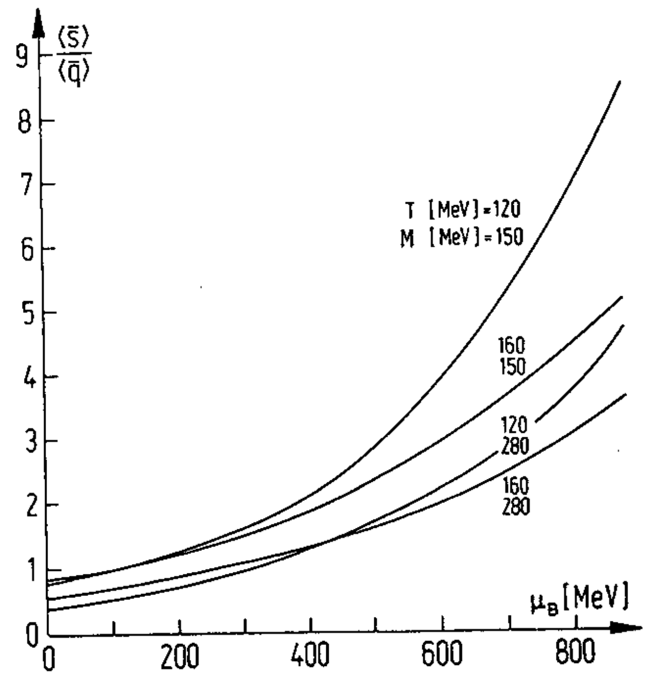


FIGURE 1 : Abundance of strange (=antistrange) quarks relative to light quark abundance as function of μ for several choices of T(=120,160 MeV) and strange quark mass (M=150,280 MeV)

We now find that there are often more $\bar{s}$ quarks than anti-quarks of each light flavour. Indeed:

$$\bar{s}/\bar{q} = \frac{1}{2}\left(\frac{M}{T}\right)^2 K_2\left(\frac{M}{T}\right) e^{\frac{\mu}{3T}} \quad . \tag{1.3}$$

This ratio is shown in Fig. 1. Thus we almost always have more $\bar{s}$ than $\bar{q}$ quarks and, in many cases of interest, $\bar{s}/\bar{q} \sim 5$. As $\mu \to 0$ there are about as many u and d quarks as there are s quarks at $T \gtrsim M = 150$ MeV.

Thus strangeness is very abundant in the plasma. The crucial aspects of the proposal to use strangeness as a tag of quark gluon plasma involve

(a) assumption of thermal *and* chemical equilibrium (see section 3);

(b) comparison between results anticipated in both hadronic phases at given T and μ, the chemical potential to be determined by other considerations (see next section and section 3).

2. STRANGE PARTICLES IN HOT NUCLEAR GAS

My intention in this section is to establish quantitatively the different channels in which the strangeness, however created in nuclear collisions, will be found. In our following analysis a tacit assumption is made that the hadronic gas phase is practically a superposition of an infinity of different hadronic gases, and all information about the interaction is hidden in the mass spectrum[3] $\tau(m^2,b)$ which describes the number of hadrons of baryon number b in a mass interval dm^2 and volume $V \sim m$. When considering strangeness-carrying particles, all we then need to include is the influence of the non-strange hadrons on the *baryon chemical potential* established by the non-strange particles. The total partition function is approximately multiplicative in these degrees of freedom:

$$\ln Z = \ln Z^{\text{nonstrange}} + \ln Z^{\text{strange}} \quad . \tag{2.1}$$

For our purposes, i.e. in order to determine the particle abundances, it is sufficient to list the strange particles separately and we find

$$\ln Z^{\text{strange}}(T,V,\lambda_s,\lambda_q) = C\{2W(x_K)[\lambda_s\lambda_q^{-1} + \lambda_s^{-1}\lambda_q] + 2[W(x_\Lambda) + 3W(x_\Sigma)][\lambda_s\lambda_q^2 + \lambda_q^{-1}\lambda_q^{-2}]\} \tag{2.2}$$

where

$$W(x_i) = \left(\frac{m_i}{T}\right)^2 K_2\left(\frac{m_i}{T}\right) \quad . \tag{2.3}$$

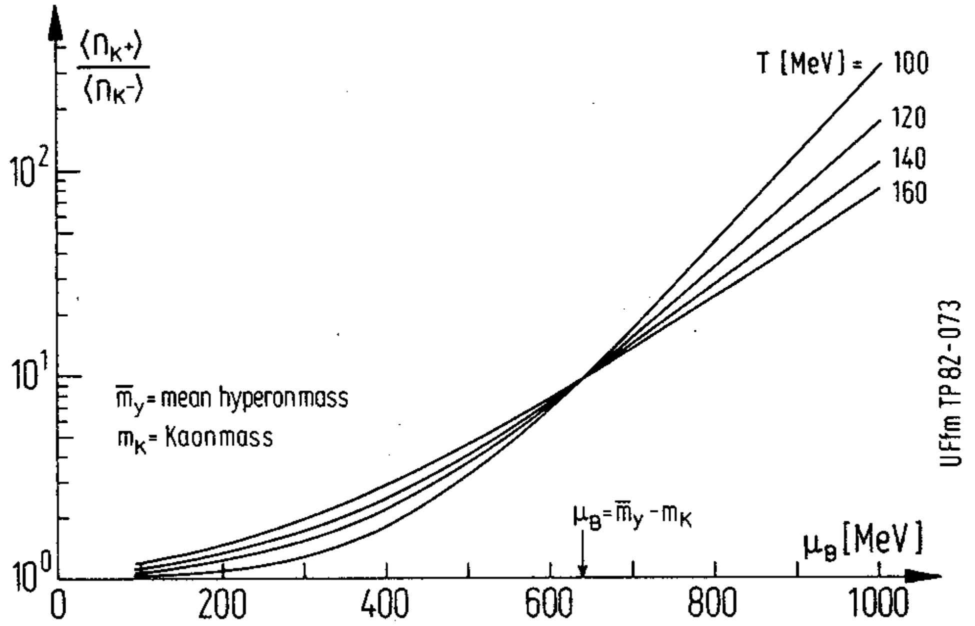


FIGURE 2 : The ratio $\langle n_{K^+}\rangle/\langle n_{K^-}\rangle = \gamma^{-2}$ as function of the baryo-chemical potential for several temperatures.

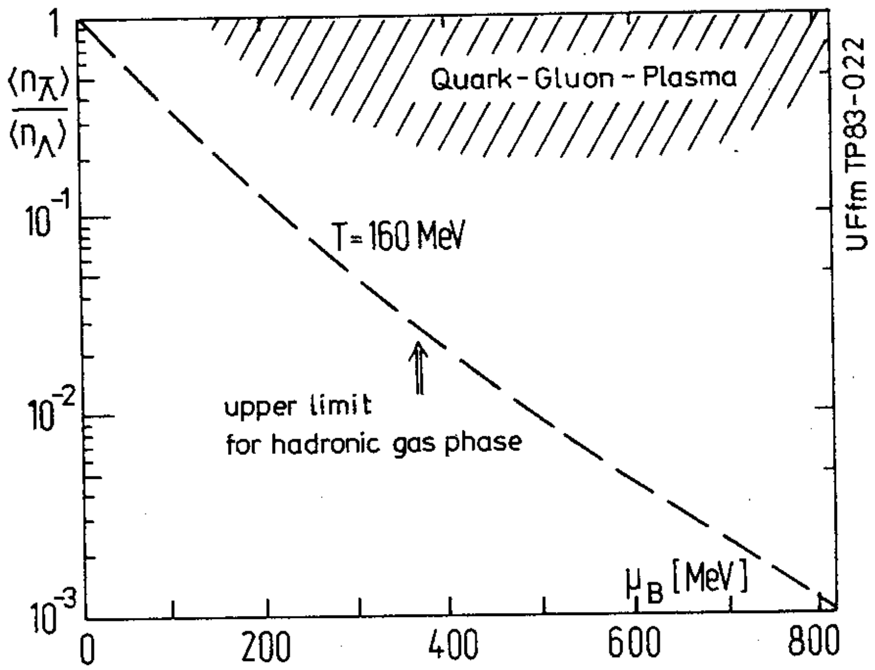


FIGURE 3 : Relative abundance of antilambdas: the actual yield from the hadronic gas limit may be still 10-100 times smaller than the statistical value shown.

We have $C = VT^3/2\pi^2$ for a fully equilibrated state. However, strangeness-creating ($x \to s + \bar{s}$) processes in hot hadronic gas may be too slow (see below) and the total abundance of strange particles may fall short of this value of C expected in *absolute strangeness chemical equilibrium*. On the other hand, strangeness-exchange cross-sections are very large (e.g., K^-p cross-section is $\sim$100 mb in the momentum range of interest), and therefore any momentarily available strangeness will always be distributed among all particles in (2.2) according to the values of the fugacities $\lambda_q = \lambda_B^{1/3}$ and λ_s. Hence we can speak of *relative strangeness chemical equilibrium*.

We neglected to write down quantum statistics corrections as well as the multistrange particles, Ξ and Ω^-, as our considerations remain valid in this simple approximation[8]. Interactions are effectively included through explicit reference to the baryon number content of the strange particles as just discussed. Non-strange hadrons influence the strange fraction by establishing the value of λ_q at the given temperature and baryon density.

The fugacities λ_s and λ_q as introduced here control the strangeness and the baryon number respectively. While λ_s counts the strange quark content, the up- and down-quark content is counted by $\lambda_q = \lambda_B^{1/3}$.

Using the partition function Eq.(2.2) we calculate for given μ_B, T, and V the mean strangeness by evaluating

$$\langle n_s - n_{\bar{s}} \rangle = \lambda_s \frac{\partial}{\partial \lambda_s} \ln Z^{\text{strange}}(T,V,\lambda_s,\lambda_q), \tag{2.4}$$

which is the difference between strange and antistrange components. This expression must be equal to zero due to the fact that the strangeness is a conserved quantum number with respect to strong interactions. From this condition we get:

$$\lambda_s = \lambda_q \left| \frac{W(x_K) + \lambda_B^{-1}[W(x_\Lambda) + 3W(x_\Sigma)]}{W(x_K) + \lambda_B[W(x_\Lambda) + 3W(x_\Sigma)]} \right|^{\frac{1}{2}} \equiv \lambda_q \gamma , \tag{2.5}$$

a result contrary to intuition: $\lambda_s \neq 1$ for a gas with total $\langle s \rangle = 0$. We notice a strong dependence of γ on the baryon number. For large μ_B the term with λ_B^{-1} will tend to zero and the term with λ_B will dominate the expression for λ_s and γ. As a consequence the particles with fugacity λ_s and strangeness s = -1 (note that by convention strange quarks s carry s = -1, while strange antiquarks $\bar{s}$ carry s = 1) are suppressed by a factor λ which is always smaller than unity. Conversely, the production of particles which carry the strangeness s = +1 will be favoured by γ^{-1}. This is the consequence of the presence of nuclear matter; for $\mu = 0$ we find $\gamma = 1$.

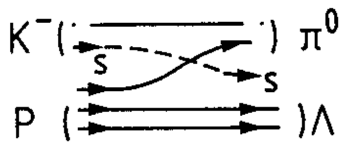
In nuclear collisions the mutual chemical equilibrium, that is, a proper distribution of strangeness among the strange hadrons, will most likely be achieved. By studying the relative yields, we can exploit this fact and eliminate the absolute normalization, C, cf., Eq. (2.2), from our considerations. We recall that the value of C is uncertain for several reasons: (i) V is unknown, (ii) C is (through space-time dependence of temperature) strongly (t,r)-dependent, and (iii) most importantly, the absolute normalization $C = VT^3/2\pi^2$ which assumes absolute chemical equilibrium, is not achieved owing to the shortness of the collision. Indeed we have [cf., Eq. (3.11) for further details and solutions]

$$\frac{dC}{dt} = A_H(1 - C^2(t)/C^2(\infty)) \qquad (2.6)$$

and the time constant $\tau_H = C(\infty)/A_H$ for strangeness production in nuclear matter can be estimated to be 10^{-22}sec.[9]. Thus C does not reach $C(\infty)$ in plasmaless nuclear collisions. If the plasma state is formed, then the relevant $C > C(\infty)$ (see below).

Now, why should we expect *relative* strangeness equilibrium to be reached faster than *absolute* strangeness equilibrium? Consider the strangeness *exchange* reaction

$K^- p \to \Lambda\pi^0$



(2.7)

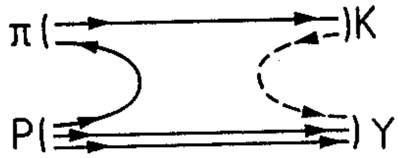
which has a cross-section of about 10mb at low energies while the $s\bar{s}$ "strangeness-creating" associate production

$pp \to p\Lambda K^+$

(2.8)

has a cross-section of less than 0.06mb, i.e., 150 times smaller. Since the latter reaction is somewhat disfavoured by phase space, consider further the reaction (Y is any hyperon = strange baryon)

$\pi p \to Y K$



(2.9)

which has a cross-section of less than 1 mb, still 10 times weaker than *one* of the s exchange channels (2.7). Consequently, I expect the relative strangeness equilibration time to be about ten times shorter than the absolute strangeness equilibration time, namely 10^{-23} sec., in hadronic matter of about twice nuclear density.

We now compute the relative strangeness abundances expected from nuclear collisions. Using Eq. (2.5) we find from Eq. (2.2) the grand canonical partition sum for zero average strangeness:

$$\ln Z_0^{\text{strange}} = C\{2W(x_K)[\gamma\lambda_K + \gamma^{-1}\lambda_{\bar K}] + 2W(x_\Lambda)[\gamma\lambda_B\lambda_\Lambda + \gamma^{-1}\lambda_B^{-1}\lambda_{\bar\Lambda}] + 6W(x_\Sigma)[\gamma\lambda_B\lambda_\Sigma + \gamma^{-1}\lambda_B^{-1}\lambda_{\bar\Sigma}]\} , \qquad (2.10)$$

where, in order to distinguish different hadrons, dummy fugacities λ_i, $i = K, \bar K, \Lambda, \bar\Lambda, \Sigma, \bar\Sigma$ have been written: the strange particle multiplicities follow then from

$$\langle n_i\rangle = \lambda_i \frac{\partial}{\partial\lambda_i} \ln Z_0^{\text{strange}} \Big|_{\lambda_i = 1} . \qquad (2.11)$$

Explicitly we find (notice that the power of γ follows the s-quark content):

$$\langle n_{K\pm}\rangle = C\,\gamma^{\mp 1}\, W(x_K) \qquad (2.12)$$

$$\langle n_{\Lambda/\Sigma^0}\rangle = C\,\gamma^{+1}\, W(x_{\Lambda/\Sigma^0}) e^{+\mu_B/T} \qquad (2.13)$$

$$\langle n_{\bar\Lambda/\bar\Sigma^0}\rangle \le C\,\gamma^{-1}\, W(x_{\bar\Lambda/\bar\Sigma^0}) e^{-\mu_B/T} . \qquad (2.14)$$

In Eq. (2.14) we have indicated that the multiplicity of antihyperons can only be built up if antibaryons are present according to their (small) phase-space. This still seems an unlikely proposition and the statistical approach may be viewed to provide an upper limit on their multiplicity.

From the above equations we can derive several very instructive conclusions[10]. In Fig. 2 we show the ratio $\langle n_{K^+}\rangle/\langle n_{K^-}\rangle = \gamma^{-2}$ as a function of the baryon-chemical potential μ_B for several temperatures that can be expected and which are seen experimentally.

We note that this particular ratio is a good measure of the baryon chemical potential in the hadronic gas phase, provided that the temperatures are approximately known. The mechanism for this process is: the strangeness exchange reaction [Eq.(2.7)] tilts to the left (K^-) or to the right (abundance $\gamma \sim K^+$), depending on the value of the baryo-chemical potential.

In Fig.3 the upper limit for the abundance of $\bar{\Lambda}$ as measured in terms of Λ-abundances is shown. Clearly visible is the substantial relative suppression of $\bar{\Lambda}$, in part caused by the baryo-chemical potential factor, Eq.(2.14), but also by the strangeness chemistry (factor γ^2) as in K^+K^- above. Indeed, the actual relative number of $\bar{\Lambda}$ will be even smaller, since Λ are in relative chemical equilibrium and $\bar{\Lambda}$ are not: the reaction $K^+\bar{p} \to \bar{\Lambda}\pi^0$, analogue to (2.7), will be suppressed by *low* $\bar{p}$ abundance. Also indicated in the Fig.3 is a rough estimate for the $\bar{\Lambda}$ production in the plasma phase, which suggests that anomalous $\bar{\Lambda}$ abundances may be an interesting feature of highly energetic nuclear collisions (cf. Section 4).

Before turning to strangeness production in the quark-gluon plasma, I would like to discuss the relative abundance of strangeness, assuming absolute chemical equilibrium, in both phases. In the hadronic gas phase we have (see Eqs. (2.2-2.5)

$$\langle n_s^{gas} \rangle = C\gamma\{2W(x_K) + 2[W(x_\Lambda) + 3W(x_\Sigma)]\lambda_q^3\} = \langle n_{\bar{s}}^{gas} \rangle \qquad (2.15a)$$

while in the quark-gluon plasma we find (see Eq.(1.1))

$$\langle n_s^{plasma} \rangle = C\, 6W(x_s) = \langle n_{\bar{s}}^{plasma} \rangle \qquad (2.15b)$$

where $x_s = M/T$ is the x-parameter of strange quarks. Hence we have

$$\frac{\langle n_s^{plasma} \rangle}{\langle n_s^{gas} \rangle} \gtrsim 3 \cdot \gamma^{-1} \frac{W(x_s)}{W(x_K) + (W(x_\Lambda) + 3W(x_\Sigma))e^{\mu/T}} \qquad (2.16)$$

where the greater sign indicates allowance for the fact that in the gas phase absolute equilibrium is much further away than in the plasma phase (perhaps factor 5 to 10). The factor of 3 in Eq.(2.16) is the colour factor - in the plasma all colours of strange quarks are permitted. The factor γ^{-1} indicates that in the hadronic gas phase the s-quark always replaces a light antiquark in a hadron. Expression (2.16) clearly shows that (except for unreasonably high values of $\mu \sim m_\Lambda$, which lead to energy densities of 10 GeV/fm^3 and more at T = 100 MeV), the ratio is larger than 3: note that $\gamma^{-1} = (\langle n_{K^+} \rangle / \langle n_{K^-} \rangle)^{1/2}$ is shown in Fig.2. Also, for $x_s = 1$, i.e. $M = T \sim 160$ MeV we have $W(x_s)/W(x_K) = 3.1$ so that $\langle n_s^{plasma} \rangle / \langle n_s^{gas} \rangle \gtrsim 10$ for small μ. Thus even for fully saturated gas-phase space the hadronic gas phase loses out substantially to the quark gluon plasma, since (a) s quarks have smaller mass than kaons, (b) in plasma, colour is not constrained locally, (c) each s-quark in the gas phase replaces a light quark.

3. STRANGENESS PRODUCTION IN THE QUARK-GLUON PLASMA

I now consider the time evolution of strange particle abundance as a function of the lifetime and excitation of the plasma state[11]. After identifying the strangeness-producing mechanisms we compute the relevant rates as functions of the energy density ('temperature') of the plasma state and compare them with those for light u and d quarks.

In lowest order in perturbative QCD, $s\bar{s}$-quark pairs can be created by annihilation of light quark-antiquark pairs (Fig. 4a) and in collisions of two gluons (Fig. 4b). The averaged total cross-sections for these processes were calculated by Cambridge[12]. For fixed invariant mass-squared $s = (k_1 + k_2)^2$, where k_i are the four momenta of the incoming particles ($w(s) = (1 - \frac{4M^2}{s})^{\frac{1}{2}}$):

$$\bar{\sigma}_{q\bar{q}\to s\bar{s}} = \frac{8\pi\alpha_s^{\ 2}}{27s}\left(1 + \frac{2M^2}{s}\right)w(s) \qquad (3.1a)$$

$$\bar{\sigma}_{gg\to s\bar{s}} = \frac{2\pi\alpha_s^{\ 2}}{3s}\left[\left(1 + \frac{4M^2}{s} + \frac{M^4}{s^2}\right)\tanh^{-1}w(s) - \left(\frac{7}{8} + \frac{31}{8}\frac{M^2}{s}\right)w(s)\right]. \qquad (3.1b)$$

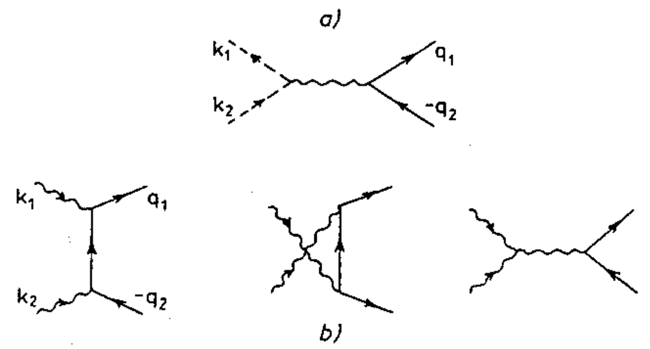


FIGURE 4 : Lowest-order QCD diagrams for $s\bar{s}$-production:
(a) $q\bar{q}\to s\bar{s}$ (b) $gg\to s\bar{s}$

For the mass of the strange quark we assume (a) the value fitted within the MIT bag model: M = 280 MeV and (b) the value found in the study of quark currents: M = 150 MeV. When discussing light quark production we use M = 15 MeV. The effective QCD coupling constant $\alpha_s = g^2/4\pi$ is an average over space- and time-like domains of momentum transfers in reactions shown in Fig. 4. We use: (a) $\alpha_s = 2.2$, the value consistent with $M = 280$ MeV in the MIT bag model, and (b) the value $\alpha_s = 0.6$, expected at the involved momentum transfer.

Given the averaged cross-sections it is easy to calculate the rate of events per unit time, summed over all final and initial states:

$$\frac{dN}{dt} = \int d^3x \int \frac{d^3k_1}{(2\pi)^3|k_1|} \sum_i \rho_i(k_1,x) \int \frac{d^3k_2}{(2\pi)^3|k_2|} \sum_i \rho_i(k_2,x) \int_{4M^2}^{\infty} ds\,\delta(s-(k_1+k_2)^2)\,k_1^{\mu}k_{2\mu}\,\bar{\sigma}(s) \tag{3.2}$$

The sum over initial states involves the discrete quantum numbers i (colour, spin, etc.) over which Eq. (3.1) was averaged. The factor $\frac{k_1 \cdot k_2}{|k_1||k_2|}$ is the relative velocity for massless particles, and we have introduced a dummy integration over s in order to facilitate the calculations. We now replace the phase space densities $\rho_i(k,x)$ by momentum distributions $f_g(k)$, $f_q(k)$, $f_{\bar{q}}(k)$ of gluons, quarks and antiquarks that can still have a parametric space-dependence, i.e. through a space dependence of temperature $T = T(x)$. The (invariant) rate per unit time and volume for the elementary processes shown in Fig. 4 is then:

$$A = \frac{dN}{dt\,d^3x} = \frac{1}{2} \int_{4M^2}^{\infty} s\,ds\;\delta(s-(k_1+k_2)^2) \int \frac{d^3k_1}{(2\pi)^3|k_1|} \int \frac{d^3k_2}{(2\pi)^3|k_2|} \cdot$$

$$\cdot\{(2\times 8)^2 f_g(k_1) f_g(k_2)\bar{\sigma}_{gg\to s\bar{s}}(s) + 2\times(2\times 3)^2 f_q(k_1) f_{\bar{q}}(k_2)\bar{\sigma}_{q\bar{q}\to s\bar{s}}(s)\}, \tag{3.3}$$

where the numerical factors count the spin, colour and isospin degrees of freedom.

We furthermore assume that in rest frame of the plasma the distribution functions f only depend on the absolute value of the momentum, $|\vec{k}| = k_0 \equiv k$; we evaluate angular integrals in Eq. (3.3):

$$A = \frac{8}{\pi^4} \int_{4M^2}^{\infty} s\,ds\;\bar{\sigma}_{gg\to s\bar{s}}\,[\int_0^{\infty} dk_1 \int_0^{\infty} dk_2\,\Theta(4k_1k_2-s)\;f_g(k_1)\;f_g(k_2)] +$$

$$+\;\frac{9}{4\pi^4} \int_{4M^2}^{\infty} s\,ds\;\bar{\sigma}_{q\bar{q}\to s\bar{s}}\,[\int_0^{\infty} dk_1 \int_0^{\infty} dk_2\,\Theta(4k_1k_2-s)\;f_q(k_1) f_{\bar{q}}(k_2)]\;, \tag{3.4}$$

where the step function Θ requires that $k_1 k_2 \geq \frac{s}{4} \geq M^2$. We now turn to the discussion of the momentum distribution and related questions. We note that the anticipated lifetime of the plasma created in nuclear collisions is of the order $6\,fm/c = 2 \cdot 10^{-23}$ sec.[13]. After this time the high internal excitation will most likely have dissipated to below the energy density required for the global restoration of the perturbative QCD vacuum state[14]. The transition between the hadronic and the quark-gluon phase is expected at an energy density of approximately 1 GeV/fm³. Under these conditions, it is possible to estimate that each perturbative quantum (light quark, gluon) in the plasma state will rescatter several times during the lifetime of the plasma. Hence the momentum distribution functions f(p) can be approximated by the statistical Bose (Fermi) distribution functions.

$$f_g(p) \approx (e^{\beta \cdot p} - 1)^{-1} \text{ (gluons)} \tag{3.5}$$

$$f_{q/\bar{q}}(p) \approx (e^{\beta \cdot p} \lambda^{\pm} + 1)^{-1}, \text{ (quarks-antiquarks)} \tag{3.6}$$

where $\beta \cdot p = \beta_0 |\vec{p}| - \vec{\beta} \cdot \vec{p}$ for massless particles, $(\beta \cdot \beta)^{-\frac{1}{2}} = T$ is the temperature-like parameter and $\lambda^{(\pm)}$ is the baryon number (antibaryon number) fugacity. In the rest frame of the plasma $\beta \cdot p = |\vec{p}|/T$. The distributions (3.5, 3.6) can only be taken seriously for $|\vec{p}|$ not much larger than T; to populate the high energy tail of the distributions too many collisions are required, for which there may not be enough time during the lifetime of the plasma. While in each individual nuclear collision the momentum distribution may vary, the ensemble of many collisions may lead to better statistical distributions.

Finally let us discuss the values of the fugacities $\lambda^{\pm}$ in Eq. (3.6). As quarks are brought into the reaction by the colliding nuclei, baryon number conservation permits to relate the baryon density ν to the fugacities by integrating Eq. (3.6) over all momenta:

$$\nu(T, \lambda^+, \lambda^-) = \frac{1}{3} \times 12 \int \frac{d^3p}{(2\pi)^3} \left[(e^{|p|/T} \lambda^+ + 1)^{-1} - (e^{|p|/T} \lambda^- + 1)^{-1} \right]. \tag{3.7}$$

The factor 1/3 takes into account the fractional baryon number of quarks. As we will show, the $gg \to q\bar{q}$ reaction time is much shorter than that for $q\bar{q} \to s\bar{s}$ production since the light quark masses are only of the order of ~15 MeV. Consequently, we may assume chemical equilibrium between q and $\bar{q}$:

$$\lambda^+ = \frac{1}{\lambda^-} = e^{-\mu_q/T}, \quad \mu = 3\mu_q \tag{3.8a}$$

$$\nu(T, \mu_q) = \frac{2}{3\pi^2} (\mu_q^3 + \mu_q (\pi T)^2) \;. \tag{3.8b}$$

As long as gluons dominate the plasma state, conditions at the phase transition, such as abundance of q and $\bar{q}$, will not matter for the $s\bar{s}$ abundances at times comparable to the lifetime of the plasma. Hence for the purpose of this study we will use the value $\mu_q = 300$ MeV in order to estimate the quark densities at given temperature. We can now return to the evaluation of the rate integrals, Eq. (3.4).

In the glue part of the rate A, Eq. (3.4), the k_1, k_2 integral can be carried out exactly by expanding the Bose function, Eq. (3.5), in a power series in $\exp(-k/T)$:

$$A_g = \frac{8}{\pi^4} T \int_{4M^2}^{\infty} ds\; s^{3/2} \bar{\sigma}_{gg\to s\bar{s}}(s) \sum_{n,n'=1} (nn')^{-\frac{1}{2}} K_1\left(\frac{(nn's)^{\frac{1}{2}}}{T}\right) . \qquad (3.9)$$

In the quark contribution an expansion of the Fermi function is not possible and the integrals must be evaluated numerically. It is found that the gluon contribution, Eq. (3.9), dominates the rate A. For $T/M \gtrsim 1$ we find:

$$A \approx A_g = \frac{7}{3\pi^2} \alpha_s^2 MT^3 e^{-2M/T} \left(1 + \frac{51}{14}\frac{T}{M} + \ldots\right). \qquad (3.10)$$

The abundance of $s\bar{s}$ pairs cannot grow forever; at some point the $s\bar{s}$-annihilation reaction will restrict the strange quark population. This loss term is proportional to the square of the density n_s of strange and antistrange quarks. With $n_s(\infty)$ being the saturation density at large times, the following differential equation determines n_s as function of time[36]:

$$\frac{dn_s}{dt} \approx A\,[1 - (n_s(t)/n_s(\infty))^2]. \qquad (3.11)$$

The solution is

$$n_s(t) = n_s(\infty) \frac{\tanh(t/\tau) + n_s(o)/n_s(\infty)}{1 + n_s(o)/n_s(\infty)\tanh(t/\tau)} , \qquad (3.12a)$$

which for $n_s(o) = 0$ becomes

$$n_s(t) = n_s(\infty)\tanh(t/\tau)\ ;\ \tau = n_s(\infty)/A \qquad (3.12b)$$

and is then a monotonically rising, saturating function, controlled by the characteristic time constant τ. In a thermally equilibrated plasma the asymptotic strangeness density, $n_s(\infty)$, is that of a relativistic Fermi gas ($\lambda = 1$):

$$n_s(\infty) = \frac{2\times 3}{2\pi^2} T M^2 \sum_{n=1}^{\infty} \frac{(-)^{n-1}}{n} K_2(nM/T) \quad , \qquad (3.13)$$

provided the volume V is large. It is found[11] that the relaxation time

$$\tau \approx \tau_g = \frac{9}{7}\left(\frac{\pi}{2}\right)^{1/2} \alpha_s^{-2}\, M^{1/2}\, T^{-3/2} e^{M/T} \left(1 + \frac{99}{56}\frac{T}{M} + \ldots\right)^{-1} \qquad (3.14)$$

is falling rapidly with increasing temperature.

We now discuss the numerical results for the rates, time constants and the expected strangeness abundance. In Fig. 5a we compare the rates for strangeness production by the processes depicted in Fig. 4 for the two different choices of parameters discussed below Eq. (3.1). The rate for $q\bar{q}\to s\bar{s}$ alone (shown separately) contributes less than 10 percent to the total rate.

In Fig. 5b we show the corresponding characteristic relaxation times toward chemical equilibrium, τ, defined in Eq. (3.12).

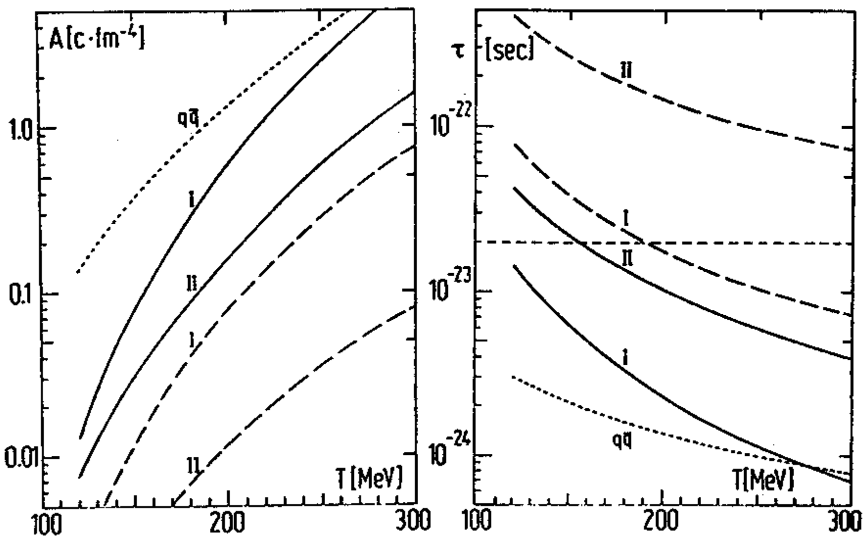


FIGURE 5 : (a) Rates A,
(b) Time constants τ as function of temperature T.
Full lines: $q\bar{q}\to s\bar{s}$ *and* $gg\to s\bar{s}$; dashed lines $q\bar{q}\to s\bar{s}$;
dotted lines $gg\to q\bar{q}$ (M=15 MeV):
Curves marked I are for α_s=2.2 and M=280 MeV,
those marked II are for α_s=0.6 and M=150 MeV.

By comparing the time constant τ with the estimated lifetime of the plasma state we find that the strangeness abundance will be chemically saturated for temperatures of 160 MeV and above, i.e. for an energy density *above 1 GeV/fm*^{3}. We note that τ is quite sensitive to the choice of the strange quark mass parameter and the coupling constant α_s which must, however, be chosen consistently. A measure of the uncertainty associated with the choice of parameters is illustrated by the difference between our results for the two parameter sets taken here.

Also included in Figs. 5 a,b are our results for gluon conversion into light quark-antiquark pairs. The shortness of τ for this process indicates that gluons and light quarks reach chemical equilibrium during the beginning stage of the plasma state, even if the quark/antiquark (i.e. baryon/meson) ratio was quite different in the prior hadronic compression phase.

The evolution of the density of strange quarks, Eq. (3.12), relative to the baryon number content of the plasma state, is shown in Fig. 6 for various temperatures. The saturation of the abundance is clearly visible for $T \geq 160$ MeV. To obtain the measurable abundance of strange quarks, the corresponding values reached after the typical lifetime of the plasma state, 2×10^{-23} sec., can be read off in Fig. 6 as a function of temperature. The strangeness abundance shows a pronounced threshold behaviour at $T \sim 120$-160 MeV.

I thus conclude, that strangeness abundance saturates in sufficiently excited quark-gluon plasma ($T > 160$ MeV, $\varepsilon > 1$ GeV/fm^3).

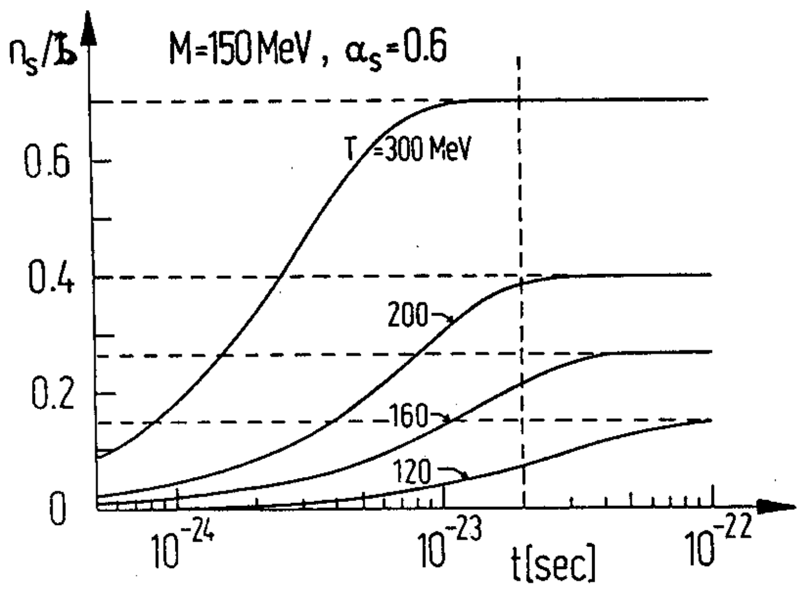


FIGURE 6 : Time-evolution of the relative strange quark to baryon number abundance in the plasma for various temperatures ($M = 150$ MeV, $\alpha_S = 0.6$).

Another remarkable fact is the high abundance of strangeness relative to baryon number in Fig. 6 - (here, baryon number was computed assuming $T \lesssim \mu_q = \frac{1}{3}\mu$).
Three facts, namely

1. high relative strangeness abundance in plasma;
2. practical saturation of available phase space in plasma, as demonstrated above;
3. s-phase in plasma up to 10 times denser than in the gas phase (see end of section 2)

suggest that the observation of strangeness is a characteristic signal of quark-gluon plasma[1]. The first point recalls (c.f.Eqs. (1.1), (1.2)) that strangeness in the quark-gluon phase is practically as abundant as the antilight quarks, $\bar{u} = \bar{d} = \bar{q}$, since both phase spaces have similar suppression factors: for $\bar{u}$, $\bar{d}$ it is the baryo-chemical potential, for $s, \bar{s}$ the mass ($M \sim \mu_q$). The third point refers to the fact that strangeness in the plasma phase is more abundant (c.f. Eq. (2.16)) than in the hadronic gas phase (even if the latter phase space is saturated) when compared at the *same* temperature and baryo-chemical potential in the phase transition region. The rationale for the comparison at fixed thermodynamic variables, rather than at fixed values of microcanonical variables such as energy density and baryon density, is outlined in the next section. We further mention here that since the abundances of anti-strangeness, both in the plasma and hadronic gas phases are equal, we also have much more anti-strangeness in the plasma as compared to the hadronic gas phase - this point can be well exploited in the discussion of the possible observables.

4. HOW TO DISCOVER THE QUARK-GLUON PLASMA

Here only the role of the strange particles in the anticipated discovery will be discussed. My intention is to show that under different possible transition scenarios, characteristic anomalous strange particle patterns emerge. Examples presented are intended to provide some guidance to future experiments and are not presented here in order to imply any particular preference for a reaction channel. I begin with a discussion of the observable quantities.

Temperature and chemical potential associated with the hot and dense phase of nuclear collision can be connected with the observed particle spectra, and, as discussed here, particle abundances. The last grand canonical variable - the volume - can be estimated from particle interferences. Thus, it is possible to use these measured variables - even if their precise values are dependent on a particular interpretational model - to uncover possible rapid changes in a particular observable. In other words, instead of considering a particular particle multiplicity as a function of the collision energy, $\sqrt{s}$, I would consider it as a function of, e.g. mean transverse momentum $\langle p_\perp \rangle$, which is a continuous function of the temperature (which is continuous across any phase transition or transformation boundary).

To avoid possible misunderstanding of what I want to say, here I consider the (difficult) observation of the width of the K^+ two-particle correlation function in momentum space as a function of the average K^+ transverse momentum obtained at given $\sqrt{s}$. Most of K^+ would originate from the plasma region, which, when it is created, is relatively small, leading to a comparatively large width. (Here I have assumed a first order phase transition with substantial increase in volume as matter changes from plasma to gas.) If, however, the plasma state were not formed, K^+ originating from the entire hot hadronic gas domain would contribute a relatively large volume which would be seen; thus the width of the two-particle correlation function would be small. Thus, first order phase transition implies a jump in the K^+ correlation width as a function of increasing $\langle p_\perp \rangle_{K^+}$, as determined in the same experiment, varying $\sqrt{s}$.

From this example emerges the general strategy of my approach: search for possible discontinuities in observables derived from discontinuous quantities (such as volume, particle abundances, etc.) as a function of quantities measured experimentally and related to thermodynamic variables always continuous at the phase boundaries: temperature, chemical potentials and pressure. This strategy, of course, can only be followed if, as stated in the first sentence of this report, approximate local thermodynamic equilibrium is also established.

Strangeness seems to be particularly useful for plasma diagnosis, because its characteristic time for chemical equilibration is of the same order of magnitude as the expected lifetime of the plasma: $T \sim 1 - 3 \cdot 10^{-23}$ sec. This means that we are dominantly creating strangeness in the zone where the plasma reaches its hottest stage - freezing over the abundance somewhat as the plasma cools down. However, the essential effect is that the strangeness abundance in the plasma is greater, by a factor of ~ 30, than that expected in the hadronic gas phase at the same values of (μ, T). Before carrying this further, let us note that in order for strangeness to disappear partially during the phase transition we must have a *slow* evolution, with time constants of $\sim 10^{-22}$ sec. But even so, we would end up with strangeness-saturated phase space in the hadronic gas phase, i.e., $\sim$ ten times more strangeness than otherwise expected. For similar reasons, i.e., in view of the rather long strangeness production time constants in the hadronic gas phase, strangeness abundance survives practically unscathed in this final part of the hadronization as well. *Facit:*

if a phase transition to the plasma state has occurred, then on return to the hadron phase there will be most likely significantly more strange particles around than there would be (at this T and μ) if the hadron gas phase had never been left.

In my opinion, the simplest observable proportional to the strange particle multiplicity is the rate of V-events from the decay of strange neutral baryons (e.g., Λ) and mesons (e.g., K_s) into two charged particles. Observations of this rate require a visual detector, e.g., a streamer chamber. To estimate the multiplicity of V-events, I reduce the total strangeness created in the collision by a factor 1/3 to select only neutral hadrons and another factor 1/2 for charged decay channels. We thus have

$$\langle n_V \rangle \approx {}^1/_6 \, \frac{\langle s \rangle + \langle \bar{s} \rangle}{\langle b \rangle} \langle b \rangle \sim \langle b \rangle / 15, \tag{5.1}$$

where I have taken $\langle s \rangle / \langle b \rangle \sim 0.2$, cf. Fig. 6. Thus for events with a large baryon number participation, we can expect to have several V's per collision, which is 100-1000 times above current observation for Ar -KCl collision at 1.8 GeV/Nuc kinetic energy[15].

Due to the high $\bar{s}$ abundance, we may further expect an enrichment of strange antibaryon abundances[1] . I would like to emphasize here $\bar{s}\bar{s}\bar{q}$ states (anticascades) created by the accidental coagulation of two $\bar{s}$ quarks helped by a gluon $\to \bar{q}$ reaction. Ultimately, the $\bar{s}\bar{s}\bar{q}$ states become $\bar{s}\bar{q}\bar{q}$, either through an $\bar{s}$ exchange reaction in the gas phase or via a weak interaction much, much later. However, half of the $\bar{s}\bar{q}\bar{q}$ states are then visible as $\bar{\Lambda}$ decays in a visual detector. This anomaly in the apparent $\bar{\Lambda}$ abundance is further enhanced by relating it to the decreased abundance of antiprotons.

Unexpected behaviour of the plasma-gas phase transition can greatly influence the channels in which strangeness is found. For example, in an extremely particle-dense plasma, the produced $s\bar{s}$ pairs may stay near to each other - if a transition occurs without any dilution of the density then I would expect a large abundance of ϕ (1020) $s\bar{s}$ mesons, easily detected through their partial decay mode (1/4%) to a $\mu^+\mu^-$ pair.

Contrary behaviour will be recorded if the plasma is cool at the phase boundary, and the transition proceeds slowly - major coagulation of strange quarks can then be expected with the formation of sss and $\bar{s}\bar{s}\bar{s}$ baryons and in general $(s)^{3n}$ clusters. Carrying this even further, supercooled plasma may become "strange" nuclear (quark) matter[16]. Again, visual detectors will be extremely successful here, showing substantial decay cascades of the same heavy fragment.

In closing this discussion I would like to give warning about the pions. From the equations of state of the plasma, we have deduced in Section 3 very high specific entropy per baryon. This entropy can only increase in the phase transition and it leads to very high pion multiplicity in nuclear collisions, probably created through pion radiation from the plasma[5] and sequential decays. Hence by relating anything to the pion multiplicity, e.g., considering K/π ratios, we dilute the signal from the plasma. Furthermore, pions are not at all characteristic for the plasma; they are simply indicating high entropy created in the collision. However, we note that the K/π ratio *can* show substantial deviations from values known in pp collisions - but the interpretations of this phenomenon will be difficult.

5. DISCUSSION AND SUMMARY

Only some selected aspects of the strangeness production in hot hadronic matter have been studied in detail -, the results are quite encouraging and suggest interesting future perspectives. In particular, it has been shown in Section 3 that strangeness abundance reaches chemical equilibrium in the plasma. The subsequent depletion of the strangeness during the plasma disintegration as well as its preferred hadronisation channels could not yet be studied in detail. However, only if the plasma disintegration is an extremely slow process, lasting of the order of 10^{-22} sec., a significant destruction of the high s-abundance created at maximal temperature reached in the collision can be anticipated: As shown in Fig. 5, the invariant rates drop quite rapidly with decreasing temperature, leading to a rapid increase of the equilibrium time constant τ - hence the strangeness abundance decouples from the absolute equilibrium by latest at the phase boundary and remains a witness of the hot collision period (see Section 4).

In Section 2 we have described how, under relative chemical equilibrium, strangeness is distributed among different hadronic channels. This being mainly controlled by the baryo-chemical potential μ, allows a measurement of μ.

There have been two messages here, that should not be missed: (1) Strangeness abundances and distribution depend greatly on the reaction channel and can greatly vary under differing experimental conditions. Thus we have a very easy and abundant diagnostic tool for the plasma state. (2) Experiments to detect the phase transition can be designed, which reflect on the possible discontinuity of the first order phase transition. They require a particular correlation between two measured quantities observed as functions of the centre of mass energy. One of these quantities varies rapidly, while the other is continuous as the phase transition is reached.

It is my opinion that experiments, in which e.g. two particle abundances are related to each other as functions of centre of mass energy, will be difficult to analyse and interpret. Here I refer in particular to the K/π ratio, which can be difficult to associate with a phase change. Similarly, measurement of the relative K^+/K^- yield, while indicative for the value of chemical potential, may carry less specific information about the plasma.

It is more appropriate to concentrate the attention on those reaction channels which will be particularly strongly populated when the quark plasma dissociates into hadrons. Here, in particular, it appears that otherwise quite rare multistrange hadrons will be enhanced, on the one side by the relative high phase space density of strangeness in the plasma, on the other side, reflecting the attractive ss-QCD interaction in the $\overline{3}_c$ and $\bar{s}s$ in 1_c channels. Hence we should search for the rise of the abundance of particles like $\Xi,\overline{\Xi},\Omega,\overline{\Omega},\phi$ and perhaps in highly strange pieces of baryonic matter, rather than in the K-channels. It seems that such experiments and others discussed in Section 4 would uniquely determine the existence of the phase transition to quark gluon plasma.

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