



THE HADRONISATION OF A QUARK-GLUON PLASMA

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Abstract

We consider two scenarios for the expansion of a quark-gluon plasma. If the evolution is slow enough, the system can remain in equilibrium throughout its entire history up to the freeze-out of a hadron gas; for a very rapid expansion, it may break up into hadrons before or at the confinement transition, without ever going through an equilibrium hadron phase.

We compare hadron production rates in the two approaches and show that for a hadronisation temperature $T \simeq 200$ MeV and baryonic chemical potential $\mu_B \lesssim 500$ MeV, their predictions essentially coincide. Present data on strange particle production lead to values in this range and hence cannot provide a distinction between the two scenarios.

Pion, nucleon and non-strange meson production seem to require a considerably lower freeze-out temperature and baryonic chemical potential. In the hadron gas picture, this is in accord with the difference in mean free path of the different hadrons in the medium; it suggests a sequential freeze-out, in which strange hadrons stop interacting earlier than non-strange hadrons.

In the quark-gluon plasma break-up, the hadronic final state fails to provide the high entropy per baryon observed in non-strange hadron production. The break-up moreover leads to a decrease of the entropy per baryon; hence it must be conceptually modified before it can be considered as a viable hadronisation mechanism.

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1. Introduction

The aim of high energy heavy ion collisions is to produce strongly interacting matter of high density. This matter will expand rapidly and eventually break up into observable hadrons. Calculations based on lattice QCD provide information about the transition of such matter from a state of deconfined coloured quarks and gluons to one of colour singlet hadrons, but they do not tell us anything about the actual evolution in which systems produced in nuclear collisions. One possible scenario is an evolution in which the system after some (early) thermalisation time remains in thermal equilibrium until the ultimately produced hadrons stop interacting. In this case, no matter what the nature of the early stages of the produced matter was, the observed hadrons will reflect the properties of the last thermal state before freeze-out, the equilibrium hadron gas. Particle production rates for this case were recently studied in detail [1]; they are given in terms of the three parameters necessary to specify the state of an equilibrium hadron gas, i.e., the temperature T , the baryon chemical potential μ_B , and the strangeness chemical potential μ_S . The requirement of vanishing overall strangeness eliminates one of these three parameters, e.g., μ_S , and then the rates are fully determined by T and μ_B . This of course presupposes the existence of a unique freeze-out point for all hadrons, which is most likely not the case. If the freeze-out is assumed to occur at the point at which the mean free path of a hadron surpasses the radius of the system, then we expect a different freeze-out for different hadrons [2]. We shall therefore consider here also the case of different freeze-out points.

An ideal hadron gas description may, however, be incorrect also for another reason. It is quite possible that the evolution proceeds too rapidly to retain equilibrium during the entire history of the system. A particularly interesting alternative is the case in which the system leaves the thermal equilibrium evolution at (or even above) the transition point to confinement [3]. In this case it is in principle possible to observe relics of the quark-gluon plasma phase in the spectra of the observed hadrons - information which is lost if the system first passes through an equilibrium hadron gas phase. A crucial feature in such a scenario are the relative weights for the different possible hadron states which can arise through quark recombination. In other words, we have to know how many s quarks combine with \bar{s} quarks to form ϕ 's, how many choose \bar{u} or \bar{d} quarks and make kaons, and how many go into strange baryons. These weights must certainly depend on the masses of the hadrons involved, on the number of possible states of the hadrons, and on the kinetic energy of the recombining partons; the latter depends on the temperature at which the break-up and recombination occur. Therefore the kinetic weight functions w_i of the different hadron species i at the transition temperature provide one possible set of such weights [4], and we shall use these here.* This means that the relative weight for recombination into hadron species a and b is just given by the ratio of the Boltzmann factors,

$$\frac{w(m_a)}{w(m_b)} = \left(\frac{d_a m_a^2}{d_b m_b^2} \right) \frac{K_2(m_a/T)}{K_2(m_b/T)}, \quad (1)$$

where d_i specifies the degeneracy (number of spin and charge states) of species i , m_i its mass, and $K_2(x)$ denotes the modified Hankel function of second order.

* For a more general discussion of recombination weights, see [5].

This choice of weights leads immediately to a problem of strangeness conservation. In the quark plasma, we assure vanishing overall strangeness by requiring an equal number of strange quarks and antiquarks; this means a vanishing strange chemical potential μ_s for quarks: $\mu_s = \mu_{\bar{s}} = 0$. This condition does not, however, provide a vanishing overall strangeness for the system of hadrons obtained through quark recombination with the relative weights given above. We have to impose strangeness conservation for the emitted hadrons in addition [4], and that reduces the number of thermodynamic parameters once more, giving us a relation between the temperature T and the chemical potential μ for nonstrange quarks at the break-up point. Hence the state of the system is now specified in terms of just one parameter.

We will now first discuss in some detail the role of strangeness conservation in the two cases, denoting the equilibrium evolution as *hadron gas* and the case of hadronisation before or at confinement as *plasma break-up*.

2. Strangeness Conservation

A. Hadron Gas

To simplify the expressions and to make the comparison to the plasma break-up easier, we here assume Boltzmann statistics; in the actual calculations, however, we shall take the correct quantum statistics into account. The partition function of the hadron gas is then given by

$$\begin{aligned} \ln Z_H(T, \mu_B, \mu_S) = & W_m + [\lambda_B + \lambda_B^{-1}] W_n \\ & + [\lambda_S + \lambda_S^{-1}] W_K \\ & + [\lambda_B \lambda_S^{-1} + \lambda_B^{-1} \lambda_S] W_Y \\ & + [\lambda_B \lambda_S^{-2} + \lambda_B^{-1} \lambda_S^2] W_{\Xi} \\ & + [\lambda_B \lambda_S^{-3} + \lambda_B^{-1} \lambda_S^3] W_{\Omega}. \end{aligned} \quad (2)$$

Here $\lambda_B \equiv \exp(\mu_B/T)$ is the fugacity for the baryon number, and $\lambda_S \equiv \exp(\mu_S/T)$ that for the strangeness. The first line corresponds to mesons and baryons of zero strangeness, the second to that of strange and antistrange mesons, the third to hyperons and antihyperons, the fourth and fifth to the Ξ 's and Ω 's, together with their antiparticles in each case. With W_i we denote the phase space factors for the various hadron states; for a fixed hadron species, they have the form

$$w_i(m_i, T) = \frac{d_i V T m_i^2}{2\pi^2} K_2\left(\frac{m_i}{T}\right), \quad (3)$$

where d_i denotes the degeneracy of hadron state i , m_i its mass, and V the volume of the system. The exact form of W_i to be used in eq.(3) depends further on the number of possible hadron states with given quantum numbers to be included here. For example, we could take into account only π, ρ, ω and ϕ as non-strange mesons, or we could include all possible states up to a certain mass. Thus we would have in the first case

$$W_m(T) = w_\pi(m_\pi, T) + w_\rho(m_\rho, T) + w_\omega(m_\omega, T) + w_\phi(m_\phi, T), \quad (4)$$

while in the second we would include further heavy mesonic resonances with their respective degeneracies. To maintain the same basis as used in [1], we shall include all resonances listed in the Particle Data Tables up to mass 2 GeV. As in [1] we also let all resonances decay according to the branching ratios listed in the Particle Data Tables, in order to obtain the actual production rates of observed hadrons (pions, kaons, etc.).

In terms of the partition function (2), we can write strangeness conservation for the hadron gas as

$$\left[\lambda_s \left(\frac{\partial \ln Z_H(T, \mu_B, \mu_s)}{\partial \lambda_s} \right) \right]_{T, \mu_B} = 0. \quad (5)$$

Eq.(5) allows us to eliminate one of the three variables T, μ_B, μ_s . In Fig.1 we show the resulting relation between μ_B and μ_s for different T ; for further details, we refer to [1]. With this relation given, we can calculate all hadron production ratios in terms of the two variables T and μ_B at freeze-out; for a number of measured strange mesons and baryons, they are tabulated in [1].

B. Plasma Break-Up

We now turn to the corresponding problem in the second scenario. The partition function for the quark plasma is given by

$$\frac{T}{V} \ln Z_P(T, \mu, \mu_s) = \frac{37\pi^2}{90} T^4 + \mu^2 T^2 + \frac{1}{2\pi^2} \mu^4 + \left(\frac{3m_s^2 T^2}{\pi^2} \right) (\lambda_s + \lambda_s^{-1}) K_2(m_s/T), \quad (6)$$

where we have included identical massless u and d quarks, and strange quarks of mass $m_s=150$ MeV; for the latter we use the Boltzmann limit. The requirement of vanishing strangeness results on the quark level in $\mu_s = 0$. If the quark system now breaks up by recombining into hadrons according to the weights (1), then it is in general not assured that e.g. the numbers of kaons, antikaons and hyperons just lead to vanishing strangeness. We have to impose this requirement in addition to setting $\mu_s = 0$ [4]. To do this, we introduce the partition function

$$\begin{aligned} \ln Z_E(T, \mu, \mu_s) \equiv & W_m + [\lambda^3 + \lambda^{-3}] W_n \\ & + [\lambda_s \lambda^{-1} + \lambda_s^{-1} \lambda] W_K \\ & + [\lambda_s \lambda^2 + \lambda_s^{-1} \lambda^{-2}] W_Y \\ & + [\lambda_s^2 \lambda_s + \lambda_s^{-2} \lambda^{-1}] W_{\Xi} \\ & + [\lambda_s^3 + \lambda_s^{-3}] W_{\Omega}. \end{aligned} \quad (7)$$

Here $\lambda \equiv \exp(\mu/T)$ is the fugacity for the non-strange quarks (we neglect the difference between u and d quarks), and λ_s that for the strange quarks. Again, the different lines correspond to mesons and baryons of zero strangeness, strange and anti-strange mesons, hyperons and antihyperons, Ξ 's and Ω 's, respectively, together with their antiparticles. In eq.(7), Z_E depends on the plasma parameters T, μ , and μ_s at break-up, and on the recombination weights W_i , which we take to have the same form as in the hadron gas case above, including the subsequent resonance decays. To discuss the strangeness or

baryon number density of the hadrons formed in the recombination, we have to relate the quark chemical potentials to those of hadrons. From eqs.(2) and (7) we have

$$\mu_S = \mu - \mu_s, \quad \mu_B = 3\mu. \quad (8)$$

In terms of these variables, the plasma requirement $\mu_s = 0$ becomes

$$\mu_S = \frac{1}{3} \mu_B. \quad (9)$$

Using eq.(8) to transform Z_E into hadron variables, we now impose vanishing overall strangeness at fixed baryon density for the emitted hadrons by requiring

$$\left[\lambda_S \left(\frac{\partial \ln Z_E(T, \mu_B, \mu_S)}{\partial \lambda_S} \right) \right]_{T, \mu_B, \mu_S = \mu_B/3} = 0. \quad (10)$$

Eq.(8) provides the indicated further relation between the temperature T and the baryonic chemical potential μ_B at break-up.

Eq.(10) can of course be also written in terms of the plasma fugacities, for which it becomes

$$\left[\lambda_s \left(\frac{\partial \ln Z_E(T, \mu, \mu_s)}{\partial \lambda_s} \right) \right]_{T, \mu, \mu_s = 0} = 0. \quad (11)$$

Using eq.(7), we obtain from this

$$(\lambda - \lambda^{-1}) W_K - [(\lambda^2 - \lambda^{-2}) W_Y + 2(\lambda - \lambda^{-1}) W_{\Xi}] = 0; \quad (12)$$

note that the number density of Ω is equal to that of $\bar{\Omega}$, so that this state drops out in eq.(12). Eq.(12) can be rewritten as

$$(\lambda - \lambda^{-1}) [W_K - (\lambda + \lambda^{-1}) W_Y - 2W_{\Xi}] = 0, \quad (13)$$

leading to $\lambda = 1$ and hence $\mu = 0$ as one solution. Since we here want to consider systems of non-vanishing baryon number density, it is not relevant for us. The remaining factor of eq.(13) leads to

$$\lambda^2 - \lambda \frac{(W_K - 2W_{\Xi})}{W_Y} + 1 = 0, \quad (14)$$

which has the solutions

$$\begin{aligned} \lambda &= \left(\frac{W_K - 2W_{\Xi}}{2W_Y} \pm \left[\left(\frac{W_K - 2W_{\Xi}}{2W_Y} \right)^2 - 1 \right]^{1/2} \right) \\ &\equiv W \pm \sqrt{W^2 - 1}. \end{aligned} \quad (15)$$

The solution with the minus sign is excluded, since $\lambda \geq 1$ for $\mu \geq 0$; hence

$$\mu = T \ln(W + \sqrt{W^2 - 1}), \quad (16)$$

with $W \equiv (W_K - 2W_\Xi)/2W_Y$, is the solution of the hadronic strangeness conservation condition for systems of non-zero baryon number density.

From eq.(16) we can deduce the break-up temperature for vanishing baryon-number density. Setting $\mu = 0$, we get $W = 1$; if we consider only the basic hadrons K, Λ, Σ and Ξ and neglect the contribution of the all higher mass meson and baryon resonances, we obtain the relation

$$w_K = 2(w_Y + w_\Xi). \quad (17)$$

Eq.(16) effectively poses at $\mu = 0$ the "remnant" condition $n_K = 2(n_Y + n_\Xi)$, with a corresponding relation for antiparticles, carried over from the situation at non-vanishing baryon density; these relations have to hold in addition to $n_K = n_{\bar{K}}, n_Y = n_{\bar{Y}}, n_\Xi = n_{\bar{\Xi}}$. Using eq.(3) and the large argument limit of $K_2(x)$ (thus assuming $m_i \gg T_i$), we get from eq.(17)

$$T \simeq \frac{(m_Y - m_K)}{\ln\{2(m_Y/m_K)^{3/2}(d_Y/d_K)\} + R}, \quad (18)$$

$$R \equiv \ln\left\{1 + \frac{1}{2} \left(\frac{m_\Xi}{m_Y}\right)^{3/2} e^{-(m_\Xi - m_\Lambda)/T_*}\right\}, \quad (19)$$

where

With $d_K = 2, d_Y = 8, d_\Xi = 4$, and the physical values for kaon and hyperon masses, this leads to $T \simeq 190$ MeV. The use of the exact Hankel function and the inclusion of different hadron states results in only slight ($\lesssim 10\%$) changes in this value; for all states up to mass 2 GeV, we get $T \simeq 200$ MeV. In this temperature range, the contribution of R is small; hence the break-up temperature at vanishing baryon number density is essentially determined by the difference between mass and degrees of freedom of kaons and hyperons.

Let us note at this point that the other extreme, the break-up of a cold baryonic plasma at $T \simeq 0$, is not immediately accessible to our considerations. We have here assumed Boltzmann statistics, and for $T \rightarrow 0$, quantum statistics certainly become relevant. In particular, cold nuclear matter containing strange baryons also requires the presence of a kaon condensate to make the overall strangeness vanish. Evidently a correct inclusion of quantum statistics becomes essential in this limit.

The general solution of the hadronic strangeness conservation equation is shown in Fig.2; it is obtained from eq.(13), using all hadronic states up to 2 GeV mass in W_K, W_Y and W_Ξ . We note in particular that over a considerable range, $T \sim \mu$, with the slope given by the weakly T -dependent function $\ln(W + \sqrt{W^2 - 1}) < 0$.

C. Hadron Gas vs. Plasma Break-Up in Strange Particle Production

Both approaches give us predictions for the ratios of produced hadrons. Where do these predictions agree, and where can we distinguish between the two scenarios? The final state is in both cases specified by a hadron-like partition function; the crucial difference between Z_H (eq.2) and Z_E (eq.7) is the plasma condition $\mu_S = \mu_B/3$ in the latter function; this relation will in general not be satisfied in the hadron gas. We see from Fig.1, however, that it is approximately true if the temperature is about 200 MeV and $\mu_B \lesssim 500$ MeV. For these parameter values the hadron gas happens to satisfy the relation obtained from strangeness conservation in the plasma, and so in this region the

two pictures are essentially indistinguishable. To illustrate that more specifically, we show in Fig.3 the ratios of the some commonly measured strange species for a hadron gas freeze-out temperature $T = 190$ MeV. As contrast we show in Fig.4 the same ratios at the freeze-out temperature $T = 120$ MeV.

The presently available data for strange particle production tend to favour the region of overlap of the two approaches; this is known from previous theoretical work [1,4] and supported by an experimental temperature determination through P_T -spectra [6]. To demonstrate it, we show in Figs. 5 and 6 the comparison of the predictions with data on $\Lambda/\bar{\Lambda}, \Xi/\bar{\Xi}, \Xi/\Lambda$, and $\Xi/\bar{\Lambda}$ [6]. This data was taken in $S - W$ collisions at central rapidity ($2.3 \leq y_{ab} \leq 3.0$), but for rather large transverse momenta. Using the experimentally determined exponential $M_T = (P_T^2 + M^2)^{1/2}$ spectra, the data can, however, be extrapolated to lower P_T in order to obtain a comparison also of Λ and Ξ as well as of their antiparticles, as shown in Fig.6 [1]. The parameters determined from these data give good agreement with the result $K_S^*/\Lambda = 1.2 \pm 0.3$ in $S - Pb$ collisions at central rapidity [7] and with $K^+/K^- = 2.3 \pm 0.5$ from $O - Au$ collisions at a somewhat lower rapidity ($1 \leq y_{ab} \leq 2$) [8]; this is shown in Fig.7. From the presently available and kinematically compatible data, the parameter values $T \simeq 200$ MeV and $\mu_B \simeq 300$ MeV emerge as hadronisation parameters for strange particle production, in accord with the conclusions of [1]. The most pressing feature to be confirmed here is that the results of [6] really persist at lower P_T . Data from [8,7] tend to give somewhat higher values for $\Lambda/\bar{\Lambda}$; this would decrease the baryonic chemical potential or increase the temperature. As it stands, the strange hadron results fall well into the range in which our two scenarios coincide (see Fig.1). Does this remain the case if we include also non-strange hadrons?

3. Production of Non-Strange Hadrons and Entropy Evolution

We now consider some results from the production of non-strange mesons and baryons. In [9,10], the ratio $\phi/(\rho + \omega)$ was found to be 4.0 ± 0.5 , measured for two different rapidity ranges and complementary P_T coverage for $O/S - U$ and $S - W$ collisions, in the dimuon channel.* From [11], we have data on the ratio h^-/p from $S - S$ interactions, giving 7.0 ± 2.0 in central collisions at central rapidity; we shall assume the negative hadrons to be pions, since the same experiment finds $K^-/\pi^- \simeq 0.06$. The value for h^-/p is supported by emulsion data for $S - Pb$ collisions, with $(h^+ + h^-)/2(h^+ - h^-) \simeq (h^-/p) + 0.5 \simeq 6$ [12]. In the hadron gas picture, the quoted ratios lead to Fig.8. We thus find that the freeze-out temperature is around $T \simeq 130$ MeV; the baryon chemical potential also shifts to lower values, around $\mu_B \simeq 200$ MeV. These results must certainly be confirmed by further experiments; but from what we know today, in an equilibrium hadron gas, the freeze-out for non-strange hadrons occurs at a noticeably lower temperature and baryon density than that for strange hadrons.

There is no unambiguous way of defining freeze-out. Usually it is assumed to occur when the mean free path of hadrons in matter has become equal to the linear dimension of the system. This gives us

$$R_F = \lambda \equiv (1/n\sigma), \quad (20)$$

* To convert this into the overall rates, we shall assume that the branching ratio of the ω into muons is the same as that into electrons.

where the mean free path is given in terms of the density of the medium

$$n = (dN/dy)/(4\pi R_F^3/3) \quad (21)$$

and the interaction cross-section σ . In eq.(21), dN/dy denotes the number of hadronic interaction partners per unit rapidity at freeze-out. From these relations, we get

$$R_F = [(3/4\pi)(dN/dy) \sigma]^{1/2} \quad (22)$$

for the freeze-out radius. The cross-sections for the interaction of strange particles are in general noticeably smaller than those for non-strange particles. The exact ratios are quite energy dependent; for a rough estimate, let us take a reduction by 0.5, in accord with $\bar{K} - p$ vs. $\pi - p$ for beam momenta around 1 GeV/c. Using eq.(22), this implies

$$R_F^{\text{str}} \simeq 0.7 R_F^{\text{n-str}}. \quad (23)$$

For spherical isentropic expansion, the freeze-out radius is inversely proportional to the freeze-out temperature ($T_F R_F^3 = \text{const.}$); hence $T_F^{\text{n-str}} \simeq 130$ MeV would imply $T_F^{\text{str}} \simeq 185$ MeV. This difference is in reasonable agreement with the temperatures obtained from our analysis of the hadron spectra.

The argument just presented does not apply directly to the annihilation of strange antibaryons in matter of non-vanishing baryon density; such reactions have on the one hand very large cross-sections, but encounter on the other hand a much lower density (dN_n/dy) of nucleons as possible annihilation partners. For 1 GeV beam momentum, the total $\bar{p} - p$ cross section is about four times bigger than the total $p - p$ cross-section. We assume the same relation for $\bar{\Lambda} - p$ in comparison to $\Lambda - p$. This increase is, however, more than compensated by a baryon density which is down by more than a factor five compared to the meson density, as we saw above. Hence the freeze-out radius of strange antibaryons is not increased by annihilation on baryons; their final interactions near freeze-out are presumably with mesons.

We now want to check the implications which the sequential nature of the hadron gas freeze-out has on the entropy of the system. Since the volume increases in the evolution, the entropy itself is not easily accessible. We can, however, study the entropy per baryon, S/B , which is just the ratio of entropy density to baryon density. Assuming that the baryon number of the system is conserved in evolution, S/B must remain constant (isentropic expansion) or increase. The entropy per baryon in the hadron gas is shown in Fig.9 for several different temperatures. We note that in the range $\mu_B \leq 400$ MeV it is considerably higher at $T = 120$ MeV than at $T = 200$ MeV. This difference is due to an increase of the entropy per baryon in the non-strange sector when we lower the temperature; the entropy per baryon of the strange sector remains nearly constant. Specifically, we find at $T = 200$ MeV about 40% of the entropy per baryon in strange particles, at 120 MeV only about 25%. At the final freeze-out ($T = 130$ MeV, $\mu_B = 200$ MeV), we obtain in the sequential hadron gas approach an overall entropy per baryon $(S/B)_H \simeq 60$; of this, about 25% comes from the strange sector, 75% from non-strange particles. Note that this high value of (S/B) is essentially due to the observed high π/p ratio; it is in accord with recently reported emulsion data [13], which lead to $(S/B) \simeq 50$ [12].

Let us now return to the plasma break-up and look at the entropy of the hadronic final state. The entropy density is obtained from eq.(7) through

$$s_E = \left(\frac{1}{V}\right) \left(\frac{\partial(T \ln Z_E)}{\partial T}\right)_{\mu_s = \mu_B/3}, \quad (24)$$

the baryon density from

$$n_E = \left(\frac{T}{V}\right) \left(\frac{\partial \ln Z_E}{\partial \mu_B}\right)_{T, \mu_s = \mu_B/3}. \quad (25)$$

Both these expressions are valid on the $T - \mu_B$ curve defined by the strangeness conservation condition (10); the ratio $(s_E/n_E) = (S/B)_E$ then gives us the entropy per baryon of the system of hadrons formed in the recombination. Its functional behaviour is shown in Fig.10. We note that at the recombination point determined from strange particle production ($T \simeq 200$ MeV, $\mu_B \simeq 300$ MeV), it gives with $(S/B)_E \simeq 15 - 20$ less than half of the experimentally observed $(S/B) \simeq 50 - 60$.

Let us compare the $(S/B)_E$ just obtained to the corresponding value in the quark-gluon plasma just before recombination. The entropy density $s(T, \mu)$ of the plasma at fixed values of T and μ , with $\mu_s = 0$, is obtained from eq.(6) as

$$s(T, \mu) = \frac{74\pi^2}{45} T^3 + 2\mu^2 T + \frac{6m_s^3}{\pi^2} \left[K_1\left(\frac{m_s}{T}\right) + \frac{4T}{m_s} K_2\left(\frac{m_s}{T}\right) \right]; \quad (26)$$

here the first term (proportional to T^3) contains contributions from gluons as well as from quarks. The baryon density n in the plasma is given by

$$n = \frac{2}{3} \left(\mu T^2 + \frac{\mu^3}{\pi^2} \right), \quad (27)$$

where $\mu = \mu_B/3$ is the chemical potential of the non-strange quarks; since the plasma contains strange quarks and antiquarks in equal numbers, they do not contribute to the baryon density. From eqns.(24) and (25), we thus get

$$(S/B)_P = \frac{x^3 \{ (37\pi^2/15) + (9/\pi^2)(m_s/T)^3 [K_1(m_s/T) + 4(T/m_s)K_2(m_s/T)] \} + 3x}{x^2 + 1/\pi^2} \quad (28)$$

for the entropy per baryon in the plasma at the moment of the break-up; here $x \equiv T/\mu$. The resulting form for the break-up parameters $T \simeq 200$ MeV, $\mu_B = 300$ MeV is included in Fig.10. We see that at the parameter values obtained from strange particle production, the plasma leads to $(S/B)_P \simeq 45$; this is of the same magnitude as the experimentally observed final overall entropy per baryon. But in particular we note that the entropy per baryon of the hadronic final state after break-up is always much lower than that of the quark-gluon plasma before break-up. The ratio of the two entropy values is shown in Fig.11. Since in any real physical process, the entropy cannot decrease, the plasma break-up approach in the form discussed so far cannot be complete.

One crucial factor which has so far been neglected in the break-up is that in an ideal quark-gluon plasma, almost half the entropy is in the gluon sector. In the recombination

process, this entropy has to go somewhere, and so we should probably not compare the entropy of the complete ideal plasma directly with that of the emitted hadrons [3,14]. To check the biggest possible effect of the gluon sector, let us just add its entropy per baryon to that of the hadrons. The behaviour resulting from this increase of $(S/B)_E$ is included in Fig.11. Although there is a clear improvement, we still fall short of fulfilling the second law of thermodynamics, by some 20 - 30 % near $T = 200$ MeV, by twice that around $T = 120$ MeV. Hence the plasma break-up scenario still needs some crucial further input before it can be considered a reasonable approach to the evolution of a quark-gluon plasma.

4. Conclusions

The production rates of hadrons provide one of our main tools for the study of the hadronisation stage in high energy heavy ion collisions; complementary information must come from correlation measurements (source size at hadronisation) and spectra (P_T distributions).

At this stage, and with all the caveats required when we compare results from different experiments, we conclude that particle production from an equilibrium hadron gas as final state seems to accommodate all observed features. At $T \simeq 200$ MeV and $\mu_B \simeq 300$ MeV, the strange particles freeze-out. Non-strange hadrons, because of their larger cross-section, continue to remain in strong interaction until $T \simeq 130$ MeV and $\mu_B \simeq 200$ MeV. The entropy per baryon at this point is about 50 - 60, in accord with experimental observations. This value is moreover comparable to that of a plasma with temperature and baryon chemical potential near those of the strange particle freeze-out, so that for an isentropic expansion, the confinement transition to an interacting hadron gas could well have occurred just above that point.

In this case, the non-strange hadrons form at $T \simeq 200$ MeV and $\mu_B \simeq 300$ MeV an interacting system with an entropy per baryon about twice that of an ideal hadron gas at the same temperature and chemical potential. It would therefore be of interest to study in more detail how interactions increase the entropy per baryon of a hadron gas. One possible mechanism for this is suggested by the statistical bootstrap model [15]. Here an exponentially growing resonance spectrum (beyond the mass limit of 2 GeV imposed in our calculations) leads to an entropy per particle diverging as the system approaches the critical point; it does so since the entropy density diverges more strongly than the particle number density [16]. Another possible mechanism appears in the bag model approach to the confinement transition at finite baryon number density [17]. In such a description, the correct high density limit is obtained only if a repulsion of whatever origin (excluded volume, Pauli blocking, ω -exchange) between baryons is included, e.g., by giving them a hard core. This leads to a reduction of the baryon number density at fixed μ_B , compared to that of ideal pointlike baryons, without affecting the mesons and hence the bulk of the entropy density. Thus $(S/B)_H$ here take on values higher than those of an ideal gas because of the effect of the interaction on the baryon density rather than on the entropy density.

A sequential freeze-out has immediate other observable consequences. It implies different freeze-out radii for kaon and pions, and this can be investigated in interferometry (HBT) studies. First results do show smaller kaon freeze-out radii [18]. It is of

interest to note that also event generator studies lead to smaller kaon freeze-out radii, precisely because of the difference in pion and kaon cross sections mentioned here [19]. The sequential freeze-out also implies different P_T spectra for strange and non-strange particles, provided these spectra do give a direct measure of the temperature. For this, however, other effects (initial state parton scattering, flow, etc.) must first be taken into account. - On the other hand, different freeze-out points for kaons and pions suggest that the much studied K/π ratio is more difficult to interpret than ratios of only strange or only non-strange particles.

The alternative evolution scheme discussed here, a break-up of the quark-gluon plasma into hadrons without ever going through an equilibrium hadron gas stage, is evidently of great interest, since it would allow the observed hadrons to transmit information from the plasma phase. In its present form, however, it does not appear tenable. It leads to a decrease of the entropy per baryon in the course of the break-up; this remains true even once the entropy of the gluons is taken into account. Perhaps also such a plasma break-up can occur in stages, although this seems much more difficult to justify than a hadronic freeze-out hierarchy based on mean free paths. But in any case, the violation of the second law of thermodynamics has to be "repaired" before a plasma break-up becomes a viable scenario.

Acknowledgements

K.R. thanks M. Gaździcki and L. Turko for useful discussions. H.S. thanks J. Zimányi for helpful comments and U. Heinz and J. Rafelski for stimulating objections. K.R. and E.S. gratefully acknowledge the financial support of the German Federal Ministry for Science and Technology (BMFT).

Figure Captions:

- 1: The strange chemical potential μ_S as function of the baryon chemical potential μ_B , as given in a hadron gas of vanishing overall strangeness, for different temperatures T . The solid line shows $\mu_S = \mu_B/3$.
- 2: The break-up temperature T as function of the baryon chemical potential μ_B , in the quark-gluon plasma break-up scenario.
- 3: The ratios of different strange particle production rates in the hadron gas approach, divided by the corresponding ratios in the plasma break-up scenario; the hadron gas freeze-out temperature is fixed at $T = 190$ MeV.
- 4: As in Fig. 3, but with the hadron gas freeze-out temperature fixed at $T = 120$ MeV.
- 5: Production ratios $\Lambda/\bar{\Lambda}$ and $\Xi/\bar{\Xi}$ for hadron gas at $T = 200$ MeV and plasma break-up, compared to data [6] at $\mu_B = 300$ MeV [1].
- 6: Same as in Fig.5, but for Ξ/Λ and $\bar{\Xi}/\bar{\Lambda}$.
- 7: Same as in Fig.5, but for K^0/Λ and K^+/K^- , with data from [7,8].
- 8: Hadron gas $T - \mu_B$ plane, with accessibility bands for $\phi/(\rho+\omega) = 4.0 \pm 0.2$ [9,10] and $\pi^-/p = 7.0 \pm 2.0$ [11].
- 9: Entropy per baryon in the hadron gas scenario, for freeze-out temperatures 0.11, 0.12, 0.13, 0.14 and 0.2 GeV.
- 10: Entropy per baryon in the plasma break-up scenario, for the quark-gluon plasma $((S/B)_P$, solid line) and for the final hadron state $((S/B)_E$, dashed line).
- 11: Ratio of the entropy per baryon in the hadronic final state after break-up to the entropy per baryon in the quark-gluon plasma (solid line); the dashed line shows the corresponding result when the gluon entropy per baryon is added to the hadronic final state.

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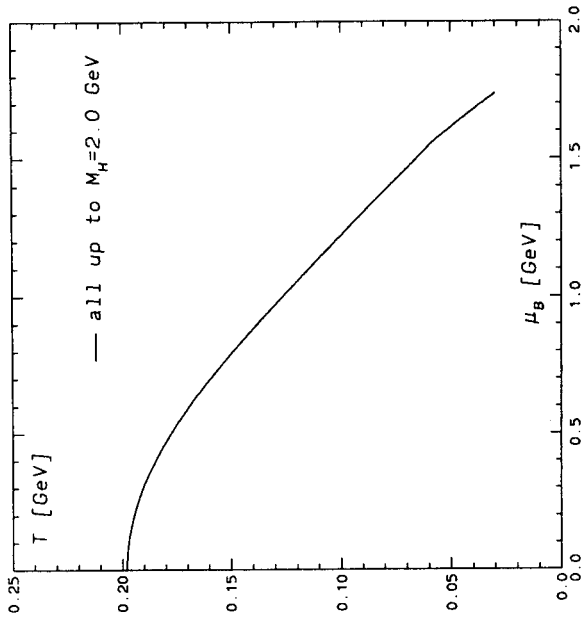


Fig. 2

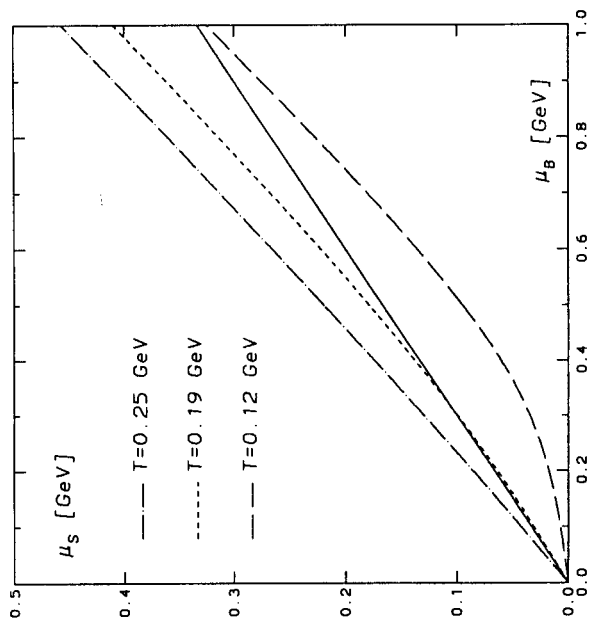


Fig. 1

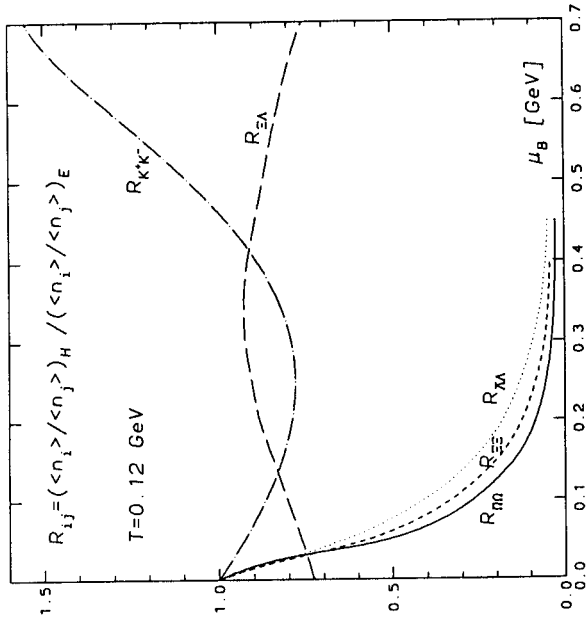


Fig. 3

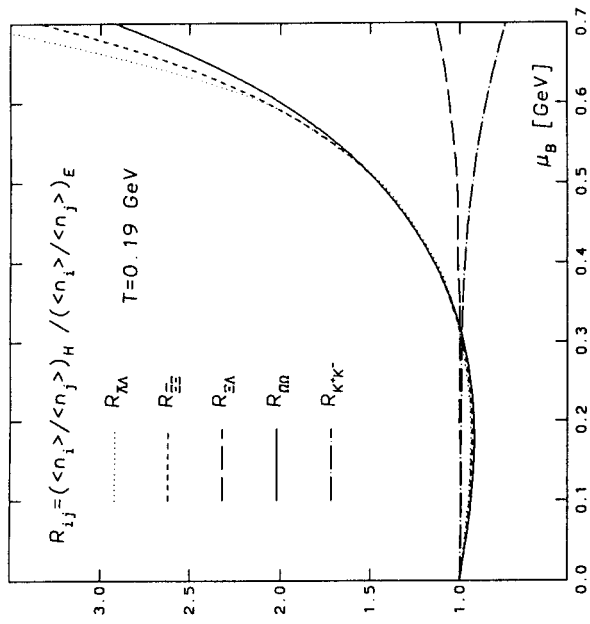


Fig. 4

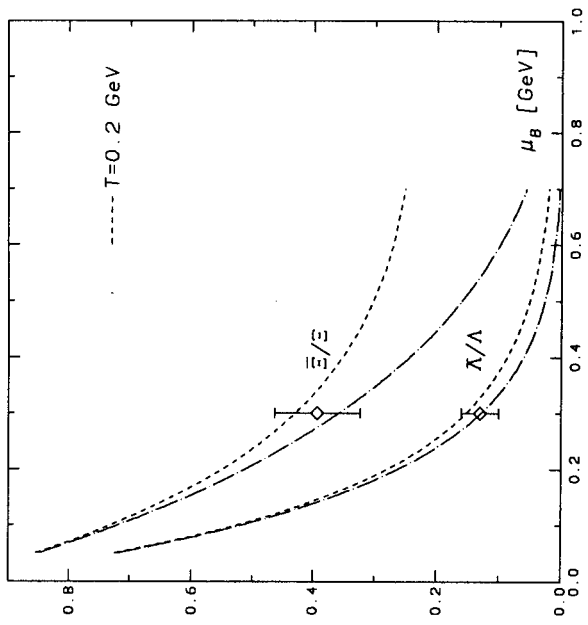


Fig. 5

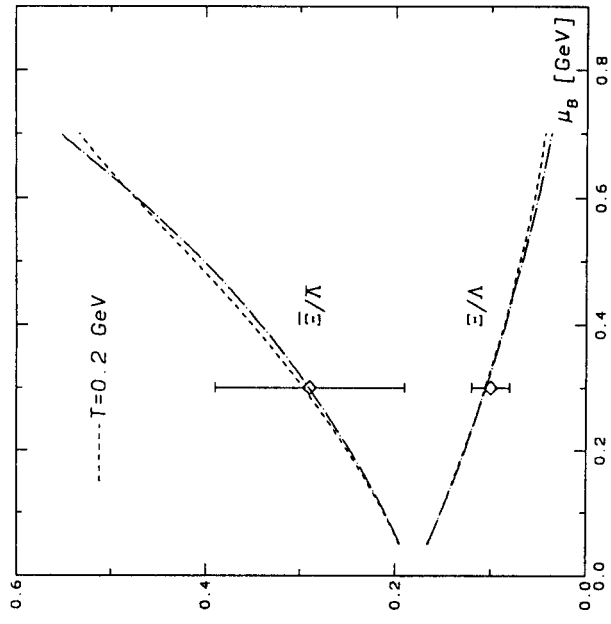


Fig. 6

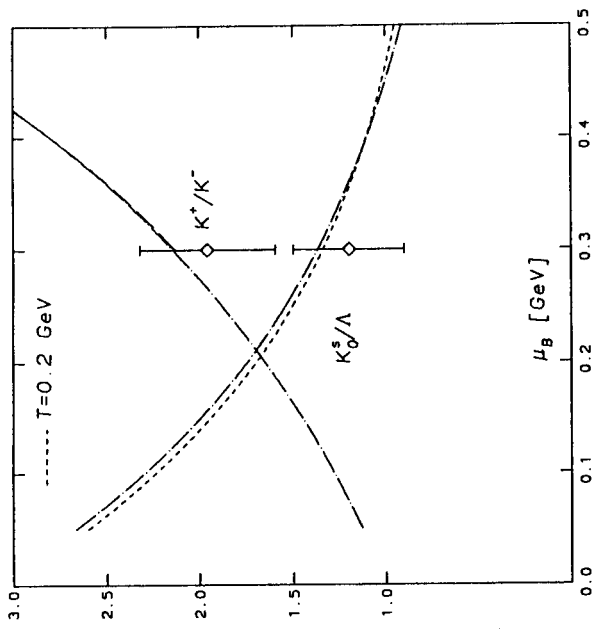


Fig. 7

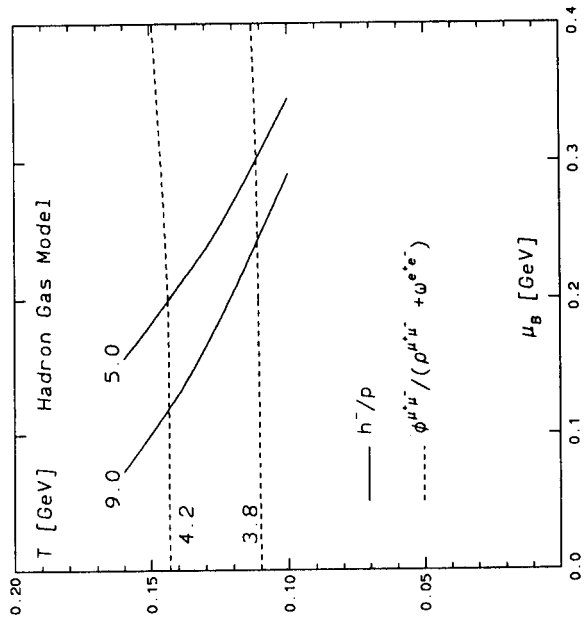


Fig. 8

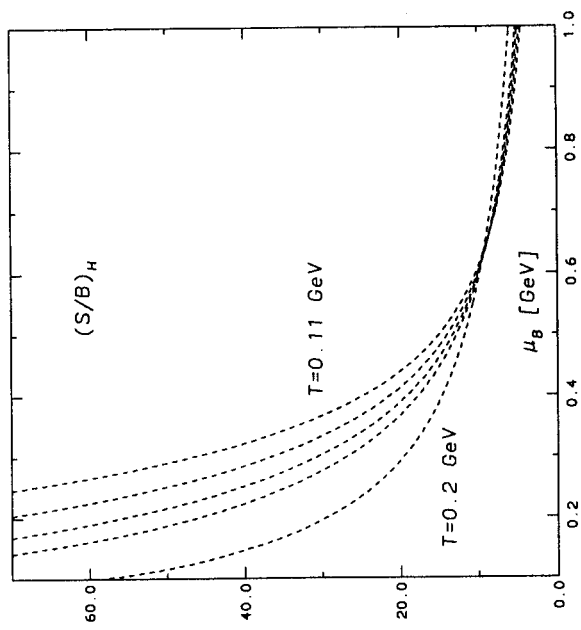


Fig. 9

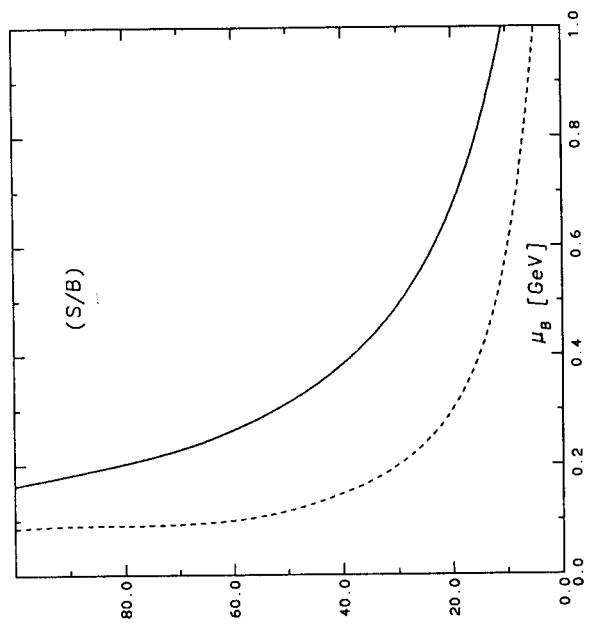


Fig. 10

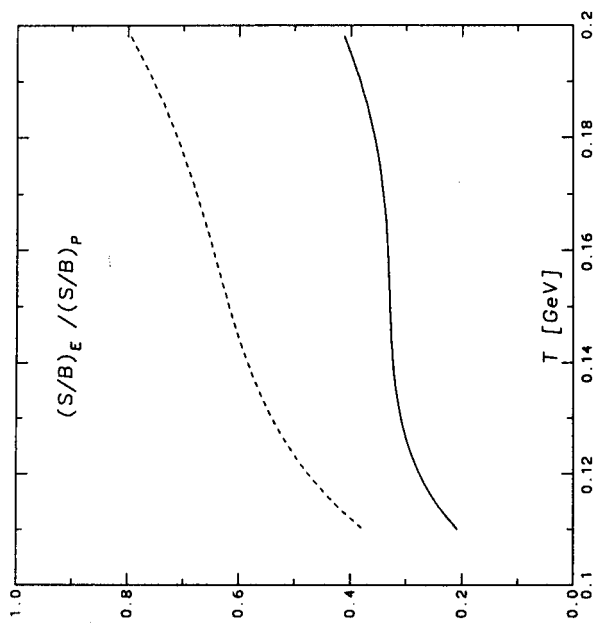


Fig.11