CERN-TH.4715/87

HOT STRANGE MATTER IN RELATIVISTIC NUCLEAR COLLISIONS*)

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A B S T R A C T

Excitation and flow of the first heavy flavour-strangeness in energetic nuclear collisions is described. Strange particles emerging from the collision are witnesses of the microscopic processes of creation of hot strange matter, especially so in the new phase-quark gluon plasma.

*) Presented at the Les Houches Meeting on "La Structure Élémentaire de la Matière", 24 March-2 April 1987.

1. INTRODUCTION

There is evidence from recent relativistic nuclear collision experiments¹⁾ that the reaction proceeds through a stage in which a hot and dense region of nuclear matter has been formed. The key issue, if the quark-gluon plasma state has been excited, will not be easy to ascertain, although this reaction mechanism is viewed as a rather compelling intermediate stage given the current QCD-based hadronic phenomenology. As it now appears, following on preliminary reports by the CERN experimental groups²⁾, there is substantial probability that the 200 GeV/A projectile Oxygen is essentially stopped by a heavy target (e.g. Pb). As function of the rapidity, the preliminary result of the NA35 collaboration shows a clear peak of negative particle multiplicity dN^-/dy around rapidity $y_{\sqrt{s}} \lesssim 2.5$, which broadens as it moves down to $y_{\sqrt{s}} \lesssim 2$ for 60 GeV/A collisions.

This and other evidence presented²⁾ suggests strongly that in central (impact parameter nearly zero) collisions, oxygen projectiles collide with the tube of matter covered by its surface, containing perhaps $\bar{A}_t = 50$ nucleons: This number is found by comparing the volume of the target tube with the nucleonic volume:

$$\bar{A}_t = \bar{V}_t / V_N = 2\pi R_t R_p^2 \frac{4\pi}{3} (1.2 \text{ fm})^3 = 1.5 A_t^{1/3} A_p^{2/3} \quad (1.1)$$

Here $\bar{A}_t$ is the nuclear content of the tube. A_t (A_p) are target (projectile) atomic numbers and $R_t(R_p) = 1.2A_t^{1/3}$ ($A_p^{1/3}$) are target (projectile) radii. The energy per nucleon in the CM frame consisting of the projectile and the tube then is

$$\sqrt{s}/A = m_N \sqrt{2\gamma_p} \frac{\sqrt{1.5(A_t/A_p)^{1/3}}}{1 + 1.5(A_t/A_p)^{1/3}} \quad (1.2)$$

where γ_p is the Lorentz factor of the projectile and the projectile rapidity is $y_p = \cosh^{-1} \gamma_p$ (6 for $\gamma_p = 200$ and $= 4.8$ for $\gamma_p = 60$). We note that there is an optimum value of A_t/A_p leading to the highest fireball energy per participating nucleon. We find this to be at $1.5(A_t/A_p)^{1/3} = 1$. Oxygen collisions at 200 GeV on $A \sim 200$ targets that the energy per nucleon in the central fireball reaches 8 GeV: should one have $A_p = 40$ available, then this value would rise to 8.5 GeV, aside from greatly increased maximum number of participating nucleons from ca 70 in the case of $A_p = 16$ to 140 participating baryons in the case of $A_p = 40$. As a final kinematic detail, we note that in the laboratory frame the rapidity of the central fireball is given by

$$y_f = \cosh^{-1} \gamma_f \quad (1.3a)$$

$$\gamma_f = \frac{A_p \gamma_p + \bar{A}_t}{\sqrt{A_p^2 + 2\gamma_p A_p \bar{A}_t + \bar{A}_t^2}} \xrightarrow{A_p = A_t} \sqrt{\frac{1 + \gamma_p}{2}} \quad (1.3b)$$

and hence for relativistic projectiles

$$y_f = \frac{1}{2} y_p - \frac{1}{2} \ln(\bar{A}_t/A_p) \quad (1.3c)$$

and we find $y_f \gtrsim 2.4$ in the particular circumstances of 200 GeV/A oxygen beam. For 60 GeV/A beam we find, similarly, $y_N \gtrsim 1.8$. We see that for the existence of central fireball, a substantial reduction of the particle rapidity from 6 (res. 4.8) to 2.5 (resp. 1.9) is required, just what is already established experimentally. Note further that in a symmetric collision consisting of e.g. N-N scattering the rapidity of the fireball would be greater than observed, with the $\ln$ -term in eq.(1.3c) vanishing.

Whatever the microscopic mechanism leading to the stopping of the projectiles in a relatively 'thin' nuclear target may be, its presence opens up the opportunity to study the properties of hot, dense, nuclear matter created in such interactions. But it is difficult to estimate the energy density ϵ reached at the moment of greatest compression. Noting that we have 8-9 GeV energy available per nucleon and that nuclear density may possibly be as compressed as five times the usual value $\rho_0 \sim 1/6 \text{ fm}^{-3}$ we arrive at $\epsilon \sim 8 \text{ GeV/fm}^3$. Such high energy density may suffice indeed for the formation of the Quark Gluon Plasma (QGP) phase which under these conditions could be as hot as 200-300 MeV, within a volume of $\sim 60 \text{ fm}^3$.

for oxygen projectiles, and exceeding 100 fm^3 for sulphur projectiles scheduled for the September/October 1987 SPS run.

With these rough reaction mechanism estimates at hand we can attempt to understand how strangeness flavour is excited in nuclear collisions. We must consider the two different scenarios, namely with and without quark gluon plasma formation. As we will see in detail in the next section, the gluonic processes in the plasma phase are highly efficient in making strange particle pairs, while the conventional associate production of strange hadrons is hindered greatly, in view of higher thresholds, lower statistical equilibrium particle density, smaller available phase space (c.f. Section 3). The nett result is that should QGP be formed, then it will contain a greatly enhanced strange particle content facilitating the formation of multiply strange hadrons. In Section 4 we consider how this will influence the abundance of otherwise rare strange hadrons, such as Ξ , and will explain how this can be used to identify the process of QGP formation and some of its properties.

2. STRANGENESS PRODUCTION IN QUARK-GLUON PLASMA

In lowest order in perturbative QCD, $(s\bar{s})$ -quark pairs can be created in collisions of two gluons and by annihilation of light quark-antiquark pairs. The gluonic production rate dominates strangeness production and leads to equilibration times comparable to the expected plasma lifetime³⁾. Given the averaged QCD reaction cross sections it is easy to calculate the rate of events per unit time summed over all final and initial states. Neglecting final state interactions (the influence of Pauli blocking is negligible) the strangeness production rate is obtained by averaging the QCD reaction cross section with the spectra of particles in the plasma: $\langle \sigma_{s\bar{s}} v_r \rangle$, with v_r being the relative velocity of colliding particles. We shall assume here that each perturbative QGP quantum (light quark, gluon) will rescatter many times during the lifetime of the plasma. Hence, we approximate the momentum distribution functions by the statistical Bose, or respectively, Fermi distribution functions, where the statistical parameters, e.g. temperature T and chemical potential μ are a priori functions of the location in the fireball ('local equilibrium').

The invariant strangeness formation rate per unit time and volume for the dominant gluonic processes is found to be³⁾

$$A_g = \frac{7}{6\pi^2} \alpha_s^2 M T^3 e^{-2M/T} \left(1 + \frac{51}{14} \frac{T}{M} + \dots \right) \quad (2.1)$$

Once strange quark-antiquark pairs have been created abundantly the $s\bar{s}$ -annihilation reaction will limit the strange quark population. Under the condition of statistical independence of the creation and annihilation process, this loss term is proportional to the square of the density ρ_s of strange and antistrange quarks. With ρ_s^∞ being the saturation density at large times, the differential equation for the evolution of ρ_s as a function of time is:

$$\frac{d\rho_s}{dt} = A[1 - (\rho_s(t)/\rho_s^\infty)^2] - \rho_s \frac{1}{V} \frac{dV}{dt} \quad (2.2)$$

The last term, the volume dilution term arises, should the volume occupied by the quark gluon plasma change in time (expansion). A and ρ_s^∞ are also time dependent quantities, through the parametric time-dependence of the temperature $T(t)$. (Both $A_g \approx A$ and ρ_s^∞ do not depend on the chemical potential of light quarks.) In Fig.2.1 the evolution of strange particle density, calculated⁴⁾ assuming $m_s = 170$ MeV, $\alpha_s = 0.6$ is shown. The simple functional relation $V \sim t$, $T \sim t^{-1/3}$, is assumed to parameterize the time dependence of T and V . Case I corresponds to a blob of plasma of size 3 fm and initial $T_0 = 250$ MeV, while case II corresponds to the initial plasma temperature being very modest, $T_0 = 180$ MeV, while the initial plasma radius being 4 fm. The energy content is almost equal in both examples. They hence represent two different conceivable scenarios of a collision of heavy nuclei with case I being favoured by current experimental results. The arrow marked 90% shows the time at which the phase space saturation reaches that level. The total strangeness is quite rapidly built up in the collision, and we can expect 80-120 strange quark pairs to be produced. The slow fall of the strangeness density is a consequence of the volume expansion, rather than strangeness annihilation. There is virtually no net strangeness annihilation $s\bar{s} \rightarrow gg$ as the plasma expands, because in the expansion process the temperature and strangeness density both drop rapidly, decoupling effectively strangeness abundance from statistical equilibrium. It is in order to illustrate this point that we have chosen the slow expansion, $V \sim t$. In fast expansion, $V \sim t^3$, the freezing out of strangeness is even more sudden. Points (a), (b), (c) and (d) indicate when the critical temperature is reached at which the chemical potential assumes the values 0, 200, 400, 600 MeV, respectively, in a particular dynamical model⁵⁾. The phase transition to the hadronic gas may occur at these points, depending on baryon content of the plasma.

As we see from Fig.2.1 the strangeness density at the transition point is expected in the interval $0.15/\text{fm}^3 < \rho_s < 0.3/\text{fm}^3$. This value indicates that clustering of two strange quarks in one hadronic volume $V_N \sim 5 \text{ fm}^3$ will be frequent, and further that the total abundance of strangeness is a rough measure of the plasma volume at the phase transition. The high particle density of strange quarks in plasma assures that when the quark matter hadronizes some of the

numerous s and $\bar{s}$ quarks may form strangeness clusters such as Ξ , Ω , and particularly important, their antiparticles $\bar{\Xi}$, $\bar{\Omega}$, instead of being bound in kaons only. Consequently, this work emphasises our expectations about these multi-strange hadrons.

Figure 2.1 : Evolution of strangeness density for two typical temporal evolution scenarios of the plasma.

3. STRANGE HADRON FORMATION OUTSIDE OF QGP (STATISTICAL MODEL)

We now study the dynamics of strangeness production in the hadronic gas phase as if no QGP were formed. The purpose of this section is to estimate the strange particle abundances expected if no quark gluon plasma has been formed in the hot fireball of the nuclear collision. When a thermally equilibrated fireball has been established, the conventional reaction for strangeness production, e.g. $N + N \rightarrow N + Y + K$ where $Y = \Lambda$ or Σ , plays almost no role, because: (1) its threshold at $\sim 670 \text{ MeV}$ energy in the center-of-mass frame is higher than in reactions between a pion and a baryon, (2) the pions are the most abundant particles except perhaps at very high values of μ_B , (3) the phasespace of two-particle final states is more favourable than that of three-body final states. Hence, the strangeness producing reactions involving a pion are dominant.

The common reaction feature of all processes is the $q\bar{q} \rightarrow s\bar{s}$ reaction, where several quarks are spectators, and a $q\bar{q}$ pair is annihilated and an $s\bar{s}$ pair is formed. The experimental value of the OZI- rule forbidden cross section for the reaction $\pi + N \rightarrow Y + K$ ($N = \text{nucleon}$), is only about 0.1 mb in the energy region of interest. Hence, the strangeness phase space saturation in HG should be a

relatively slow process since the $s\bar{s}$ forming collision time of a pion with a nucleon, at $\rho \sim 3\rho = \frac{1}{2} \frac{1}{\text{fm}^3}$ is

$$\tau \sim (\sigma v \rho)^{-1} \sim 10^{-21} \text{ s} \quad (3.1)$$

It is also clear that while the direct pair production of multiple strangeness carrying (anti-) baryons is strongly suppressed in nuclear collisions, these particles may be built up by exothermic strange quark exchange reactions. Experimentally, one finds for the strangeness exchange reaction $\bar{K} + N \rightarrow Y + \pi$, a cross section roughly ten times larger than the production cross section. This means that strangeness is much faster redistributed among the strange particle family than produced. Fast redistribution of strange s-quarks among different hadrons means that the relative chemical equilibrium of s-quarks is reached during the conceivable lifetime of the fireball, unless the fireball lives 10^{-21} sec, the absolute equilibrium. i.e. saturation of the phase space is not achieved.

Similar to eq.(2.2) we find a (now quite complex) family of evolution equations for the densities of the various particles in question. In actual calculations we used^(6,5) parameterisation of measured cross sections to calculate average cross sections entering the (nonlinear) master equations which describes the evolution of the chemical composition of the hot hadronic matter phase. The initial condition for the strange particle densities at a given 'initial' time t_0 were all set to zero, thus neglecting the contributions of strangeness production reactions in the first moments of the collision before thermal equilibrium is reached. Further below in Section 4 we will use the same master equations in order to describe the break up of QGP and the subsequent evolution particle: the hadronic reaction then starts with the strangeness abundance that has been built up in the quark-gluon plasma, as further modified in the phase transition back to hadronic matter. For the required antinucleon density at $t=t_0 \sim 10^{-24}$ sec we use the chemical equilibrium value, being aware of the fact that the assumption of baryochemical equilibrium for antinucleons is largely overestimating the possible abundance. This leads to an overestimate of the abundance of strange antibaryons generated by hadronic interactions in the absence of quark gluon plasma. As the abundance of strange antibaryons we find is extremely small anyway, our approach is a 'conservative' estimate of the background of strange antibaryons expected, should no plasma be formed.

The time evolution of the total strangeness abundance ρ_s in hadronic gas is displayed in Fig.3.1 (due to exact strangeness conservation, ρ_s is equal to the total antistrangeness density $\rho_{\bar{s}}$) for $T = 160$ MeV and two values of the baryon chemical potential, $\mu_b = 0$ and $\mu_b = 450$ MeV. When looking at both curves, we see that after a typical break-up time of the hadronic fireball of about 5×10^{-23} s,

the strangeness abundance predicted for the hadronic phase is about a factor of three smaller than its equilibrium value. Individual particles with higher strangeness content such as $\bar{\Omega}$ (see below) will be still more distant from equilibrium distribution: the density of anti-omegas as well as omegas is in HG a factor of about 10^3 lower than their equilibrium values for a breakup time between 10^{-23} and 10^{-22} seconds. This is true for baryonless as well as for baryon rich hadronic gas. Anticascades ($\bar{\Xi}$) are suppressed by a factor of about 10^2 with respect to their equilibrium values. We emphasise that in more realistic calculations we expect to find an even lower density of strange antibaryons created outside of QGP.

Figure 3.1 : Time evolution of strangeness density in hot hadronic matter at fixed temperature $T = 160$ MeV and baryochemical potentials $\mu_B = 0$ and $\mu_B = 450$ MeV.

For this reason, i.e. due to their effective absence in HQ, strange antibaryons appear as the most promising signal for the temporary existence of a deconfined quark-gluon phase during the course of the nuclear collision given the very large abundance of strangeness expected to be created in such experiments.

4. STRANGE ANTIBARYON FORMATION FROM QUARK-GLUON PLASMA

During the process of QGP hadronisation substantial fragmentation of gluons in particular is required: note that if quarks and antiquarks would just simply recombine⁷⁾ into mesons, mainly pions, there would be only half as many pions afterwards as there were quarks and antiquarks before: Each pion carries about 4

units of entropy⁸⁾, while each quark and antiquark carries 4.2 units of entropy. The entropy. Hence in the recombination half of the entropy is lost, while one also ignores another factor of two in entropy when gluons are not accounted for. Thus in a recombination model⁷⁾ nearly 3/4 of entropy disappears. But we know that the total entropy should increase during the phase transition in view of the second law of thermodynamics. In order to appreciate what went wrong, we recall that the usual chemical reactions involve nonrelativistic particles, whose entropy per particle increases with temperature and also possibly on account of internal excitations. The entropy can, therefore, be made to increase while the number of particles decreases due to recombination. Such an effect is not possible for relativistic particles. The large amount of entropy residing in the quark-gluon plasma can only be disposed of by generating, ultimately, a sufficiently large number of pions⁸⁾.

There exists a simple mechanism permitting explicit implementation of the fragmentation process: recall that a quark-antiquark pair will only in one case out of nine be found in a colour singlet state that can form a single meson. In the other eight cases the quark-antiquark pair, forming a colour octet state, must first radiate a gluon in order to be able to recombine into a meson. Similarly, gluons will not, in general, simply disappear into the vacuum during the hadronization process in view of their colour octet nature, but will e.g. fragment into a quark-antiquark pair. We find that in order to conserve entropy during the hadronization process, every gluon and about one third of the quarks must fragment before coalescing into mesons. The fragmentation processes must be correctly weighted with regard to their flavour content. Although it is not entirely clear whether the concepts developed⁹⁾ for quark jet fragmentation are valid in the environment of the hadronizing quark-gluon plasma, we shall here adopt these concepts as a first guideline. The relative probability of glue fragmentation into a quark pair of mass m_i is controlled by the parameter f_i . For light quarks we can take $m_u = m_d = 0$ leading to $f_u = f_d \sim 0.425$, $f_s \sim 0.15$ for $m_s \sim 170$ MeV. The number of quark-antiquark pairs of each flavour is now obtained by multiplying the number of fragmenting gluons with the probability f_i for producing the considered flavour.

For the description of the strangeness evolution and strange particle excitation during the spatial expansion of the fireball in its hadronization stage some further detailed model assumptions are required: there are two mechanisms that lead to a change in volume occupied by the fragmenting QGP in the mixed phase. The volume may change due to a hydrodynamical expansion at nearly constant entropy. Alternatively the volumes can change due to exchange of particles, i.e. entropy, with the hadronic gas phase. Simultaneously, the volume of the hadronic phase continuously increases, due to the accumulation of hadronic particles as well as due to (hydrodynamical) volume expansion. Note also that a large volume jump

occurs between QGP and hadronic phases as suggested by the high entropy density of the plasma as compared to the hadronic gas phase. Consequently, the master equations describing the mixed phase are quite complicated.

We are mainly interested in the fate of strange quarks. The number of strange quarks in the plasma phase changes because of two reasons: on the one hand, strangeness producing reactions make it approach the equilibrium abundance (c.f. Section 2). Furthermore, the number of strange quarks in the plasma phase changes due to their conversion into hadrons. In the hadronic phase the number of strange particles will further change through standard hadronic reactions and this part of the reaction can be treated as described in Section 3. We now summarise the results⁵⁾ of calculations involving, in particular, strange and nonstrange baryons and antibaryons in addition to pions and kaons. Note that these calculations are based on the following physical picture of the evolution of the fireball: initially a quark-gluon plasma is formed, which expands and cools until it reaches the phase transition point towards the hadronic gas phase. Temperature and baryochemical potential are then kept constant and the plasma is gradually converted into hadronic gas (mixed phase). During this process the volume of the mixed phase grows according to the requirement of entropy conservation. The calculation is stopped when the quark-gluon plasma is fully converted into hadrons. When plasma and hadronic phase coexist, complex master equations for the mixed phase are integrated in the hadronization process; mesons and baryons are assumed to form according to the fragmentation-recombination scenario - gluon fragmentation ensures approximate entropy conservation by providing the additional partons required to fill the larger hadronic gas volume with mesons and baryons.

The resulting particle abundances are presented in Fig.4.1. Parts (a) and (b) refer to abundances of hadrons containing s and $\bar{s}$ quarks, respectively. The abundances are presented as found at the time when the plasma phase has been completely converted to the hadronic gas phase and the particle density is already quite small. The curves are given as function of the values of the thermodynamical variables T and μ_b at the point where the critical curve is crossed. In the calculation a constant T/μ_b ratio was assumed. The dashed curves are the equilibrium hadronic gas abundances which, as we have discussed in the case of strange antibaryons, are gross overestimates of actual expectations for the background. As we see from the presented results, there is a very significant overabundance as compared to the hadronic gas phase space of strange antibaryons found upon hadronisation of the QGP phase. For reasons of experimental practicability it is further important that the densities of strange antibaryons $\bar{Y}$, $\bar{\Xi}$ are in the range 10^{-3} to $10^{-2}/\text{fm}^3$ at the time of hadronic break-up. Given the large fireball volume of more than a thousand fm^3 at this time, a significant number of these particles

are expected to be produced in each single reaction event leading to formation of quark gluon plasma.

Figure 4.1 : Strange particle abundances assuming fast volume expansion ($V \sim t^3$) and recombination with gluon fragmentation: (a) s - and (b) $\bar{s}$ -quark carrying hadrons at the end of coexistence of quark gluon plasma and hadronic gas phase, as a function of critical values of T and μ_B . Dashed: equilibrium abundances in the hadronic gas phase at the same values of T and μ_B .

The synopsis of the results presented here confirm the particular suitability of the global abundance of strange antibaryons for diagnosis and study of the quark gluon plasma state. However, one can also consider the direct effect of the large s , $\bar{s}$ density in the early plasma before hadronisation, and in particular the possibility of strange particle radiation from such a hot plasma state¹⁰⁾. Two channels of particle production may be considered. In the first primordial particle emission mechanism a fast quark (or coloured diquark) from the plasma impinges on the boundary between the plasma perturbative vacuum and the true vacuum. In the associated color string breaking process⁹⁾, at least one quark-antiquark pair is formed. We will refer to this process as a 'microjet' (not to be confused with the minijet processes). Second, in the recombination approach, several constituents of the plasma, clustered into colorless objects, penetrate the surface and hadronize. This latter process is similar to the model of the

phase transformation of the plasma to the hadronic gas at the later stages of the plasma life which we have just considered. However, this recombination radiation is, in detail, different in that we may not allow gluon fragmentation. The equilibrium quark abundances in plasma contain the effect of continuous gluon and quark fragmentations and recombinations. Only at the point of phase transition is equilibrium lost and additional particle and entropy generating microscopic processes such as gluon fragmentation become necessary. We expect that the microjet picture will dominate the spectra at medium to high $E_{\perp} \sim (2-4)m$.

In Fig.4.2 we compare for the $\bar{\Xi}/(\bar{Y}/2)$ ratio the results¹⁰⁾ of the microjet and recombination mechanisms to global abundance ratios i.e. (including low $p_{\perp}$), c.f. Fig.4.1. The surprise is that in both early emission processes (microjet, recombination) populating medium to high $E_{\perp}$ portions of the spectra we are led to expect more anticascades than antilambdas (since $\bar{\Sigma}^0 \rightarrow \bar{\Lambda} + \gamma$ we included $\bar{\Sigma}^0$ into the 'lambda' abundance: the visible hyperon decays are just half of the total abundance).

Figure 4.2 : $\bar{\Xi}/(\bar{Y}/2)$ particle ratio as a function of $\mu_b/3T$ for the microjet, recombination models of particle spectra at medium to high $E_{\perp}$. Dashed: global ratio.

But there is some difference which is difficult to differentiate between the two early emission processes, unless the value of μ_b is known to some accuracy. Although experimental observation of the ratio might appear difficult, longitudinal polarisation method¹¹⁾ can be applied in visual experiments in e.g. CERN

experiment NA35, while effective decay length of a Ξ indicates it as a descendant of a Ξ . This method may be applied in the CERN experiment WA85. Electronic particle track reconstruction will be attempted in CERN experiment NA36 for multiply strange, cascading decay particles.

The source of all these surprising results about strange hadrons from QGP can be traced back to the fact that strange quark-pair production in the plasma phase proceeds at a sufficiently fast rate to permit statistical equilibrium abundance to be established in times shorter than 10 fm/c. This is due to the abundant presence of gluonic excitations, allowing for quark-pair production in gluon-gluon collisions, c.f. Section 2. In a sense, therefore, abundant strange antibaryon production is indicative of an environment in which gluon collision processes are an essential element of the reaction picture.

In conclusion, we have shown on the basis of a dynamical theoretical model of strange particle production in nuclear collisions that normally rare strange antibaryon particles with multiple strangeness content provide a very promising experimental signal in the search of the quark-gluon plasma. Abundant strangeness production is indicative of the presence of gluon excitations, a characteristic property of the deconfined QCD phase.

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