Strange Particle Abundance in QGP Formed in 200 GeV A Nuclear Collisions

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Abstract

We investigate the relative abundance of strange particles produced in nuclear collisions at the SPS energies (\(\sim 9\) GeV A in CM frame) assuming that the central reaction fireball consists of quark-gluon plasma. We show that the total strangeness yield observed in Sulphur induced reactions is compatible with this picture.

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1. Introduction: Abundant strangeness production in relativistic nuclear collisions is suggestive of the quark-gluon plasma (QGP) formation [1]. But the relative abundances of strange and multi-strange baryons and anti-baryons have been suggested as the better signatures discriminating against other possible forms of highly excited hadronic matter. These relatively complex composite particles are more sensitive to the environment from which they emerge and show pronounced anomalies in their expected yields [2] as compared to experience based on N–N reactions. However, their systematic measurement constitutes a formidable experimental challenge and only today experimental results on their (relative) abundances and spectra are being reported [3, 4, 5, 6, 7] for S–A (A = S, Ag, W, Pb) collision systems.

It was immediately recognized that this data lend itself naturally to an interpretation in terms of local thermal equilibrium fireball model and the resulting thermal parameters are suggestive of a source consisting of a deconfined QGP fireball [8]. This hypothesis was further supported by the observation that details of the produced particle multiplicity point to a high entropy primordial phase [9]. Such a phase could not be brought into consistence with the properties of a conventional hadronic gas state made of individual hadrons. Rejection of confined, hadronic gas (HG) phase became even more compelling after a more complete analysis of (revised) data was completed allowing also for resonance disintegration [10, 11]. The formation of the primordial high entropy phase in turn implies that the relative yield of strange to non-strange particles will be diluted in the final hadronic state by the necessity to generate a high particle multiplicity associate with the high entropy content of the source. Hence the question ensues if indeed strangeness multiplicity enhancement is visible should the hadronic multiplicity increase, as is expected for a QGP central fireball. We will consider how the theoretical expectations about the relative abundance of strange to non-strange particles produced from the thermal QGP phase compare to the results of nuclear collisions for S–S, S–Ag [8] reaction systems and comment on S–W [4] and S–Pb [7] results. We also shall extrapolate our findings to the forthcoming Pb–Pb collision experiments.

We find good agreement between experiment and the data, once one allows for the presence of the strangeness occupancy factor $0.5 < \gamma_s < 1$ [8], where the value of $\gamma_s$ is believed to depend on the size and life span of the QGP region formed in the reaction considered [12]. The yield we find is consistently higher than one has found in conventional N–N reactions [10], and we predict a still greater enhancement for the Pb–Pb case. Thus though the global strangeness enhancement seen at one collision energy cannot be alone seen as convincing evidence for the formation of QGP state, we find that the magnitude and systematic behavior seen at 200 GeV as function of the target mass fits well into the picture of QGP fireball developed to account for the anomalies found in (multi-)strange (anti-)baryon abundances.

2. Reaction dynamics: Implicit in the physical picture employed here (see Ref. [11] for more details) is the formation of a central and hot matter fireball in the nuclear collisions at 200 GeV A. This reaction picture presumes that a fraction of energy and flavor (baryon number) content of the central rapidity region is able to thermalize, characterized by the temperature $T$ and the chemical fugacities $\lambda_i$ of the different conserved quark flavors $i = u, d, s$. This local thermal equilibrium hypothesis has been inferred from the global evidence on particle spectra and an extensive analysis [14] which included diverse distortions of the spectra caused by remaining collective longitudinal flow of unstopped matter, and
resonance decays. For example we note that the WA80 \cite{13} collaboration has presented data regarding $\pi^0$ and $\eta$ temperatures over a wide transverse mass range of $0.4 < m_\perp < 3$ GeV. In the domain $m_\perp > 1.5$ GeV where the corrections arising from disintegrations of hadronic resonances are small one can infer a source temperature of about $240 \pm 10$ MeV, very similar to what is seen for the strange particles produced at such high transverse mass \cite{4}.

Given these high temperatures we can subsume the following: early in the collision there was a maximum compression point at which the conversion of energy from relative motion into the thermal energy content of the central fireball stopped either because the kinetic pressure of colliding matter has being exceeded by the internal pressure of the fireball at the kinetic energy of the collision, or because the smallness of the target/projectile the collision reaction was ended. Evidence for longitudinal flow in rapidity hadron spectra points to this latter case in the S–A reactions at 200 GeV A energy. Following the collision compression and the ensuing rapid thermalization the matter develops collective rarefaction wave which absorbs progressively greater fraction of the thermal energy, and the system cools down till the final state particles freeze-out. They are propelled by the rarefaction wave and hence their spectral temperature is blue-shifted by the collective flow velocity. Therefore the temperature observed in the $m_\perp$ particle spectra is nearly the same that would be recorded directly from the initially highly compressed state.

Thus whichever the exact dynamical history may be, \textit{e.g.} surface evaporation or rarefaction followed by low density freeze-out or combination of both, the particle spectra carry the information about the temperature at which the heating due to the dynamical squeeze of colliding nuclei stops. The experimental evidence derived from strange particle spectra is that the initial state formed in S–A collisions at 200 GeV A reaches $T \simeq 232 \pm 3$ MeV \cite{4}, while in S–S case it is $T \simeq 194 \pm 15$ MeV \cite{6}. It is not at all self-evident that these results are consistent with the nuclear reaction process when QGP equations of state are invoked. The condition of the initial state formed in the collision is controlled by the stopping of energy and baryon number in the central fireball. We introduce coefficients $\eta_E, \eta_B$ which describe the fraction of energy and baryon number deposited in the thermalized fireball. A typical value arising from global characteristics of the reaction data is $\eta_E, \eta_B = 60\%$ of the CM available energy and the total baryon number content of S–Pb reaction system consisting of the tube of projectile matter overlapping with the tube of matter in the target \cite{15}. The implicit hypothesis that $\eta_E/\eta_B \simeq 1$ has as important consequence that the thermal state energy per baryon is equal to the initial energy per baryon in the collision system. We use the following parameters:

| System     | Momentum | $\sqrt{s_{NN}}$ (GeV) | $\gamma_s$ |
|------------|----------|----------------------|------------|
| S–S/Ag     | 200 GeV A| 9.7/9.4              | 0.56 ± 0.2 |
| S–W/Pb     | 200 GeV A| 8.8                  | 0.8 ± 0.1  |
| Pb–Pb      | 157 GeV A| 8.6                  | 1 (?)      |

The value of strangeness occupancy factor $\gamma_s$ is not overly important (we show above our final results) in our later study of equations of state and the evolution of the compressed
state. Since, to a good approximation, $\gamma_s$ enters as a multiplicative factor in the strangeness energy density (typically 20% of all) a reduction of $\gamma_s$ by 50% has similar influence on $E/B$ as has the uncertainty related to the condition $\eta_E \simeq \eta_B$. We note that when computing the collision energy of the Pb–Pb system we have assumed that the projectile momentum per nucleon will be scaled by the factor $Z/A$ (and correcting for nuclear binding) down from the $p-p$ 400 GeV/c nominal CERN-SPS value.

3. Chemical conditions and deconfinement: Studying the relative strange particle abundances it is possible to determine precisely the values of quark fugacities $\lambda$. Full analysis of the S–S [17] and S–W data [10] obtained at 200 GeV A has shown that the strange quark fugacity $\lambda_s \simeq 1$, which is not the case for the lower energy AGS results obtained with 14 GeV A projectiles [18]. The recently reported $\Omega/\Omega$ result [3] has provided another independent indication that $\lambda_s \simeq 1.0 \pm 0.1$ [10]. $\lambda_s \sim 1$ is natural for a directly disintegrating QGP phase where the symmetry between the $s$ and $\bar{s}$ quarks is reflected naturally in the value of $\lambda_s = 1$. Since this value is observed in the final state this suggests that the hadronization of the QGP phase occurs in particular conditions, given that in general the HG phase does not maintain this value of $\lambda_s$. In the HG phase, whatever the equation of state, $\lambda_s = 1$ is an exceptional condition. At final baryon number (viz. $\lambda_d$ and $\lambda_u \neq 1$) the strangeness conservation constraint requires that the number of strange and anti-strange valance quarks bound in final state hadrons are equal. This is in general incompatible with $\lambda_s = 1$. We now turn to the light quark fugacities: as there is only a slight asymmetry in the number of $u$ and $d$ quarks in the heavy nuclei used in experiments it is thus convenient to introduce the quark fugacity $\lambda_q^2 = \lambda_u \lambda_d$ and to confine the asymmetry between the number of neutrons and protons to the parameter $\delta \lambda \leq 0.03$ with $\lambda_d/\lambda_u = (1 + \delta \lambda)^2$ [10]. Relative abundances of (strange) particles can be used to fix the value of the quark fugacity to considerable precision. For S–W collisions $\lambda_q = 1.49 \pm 0.05$ is found [10], while for S–S collisions at central rapidity, $\lambda_q \simeq 1.36$ [17] best describes the data.

It is important to note that the presence of a unique value of $\lambda_q$ determining all particle abundances for each collision system has considerable implications reaching beyond the observations regarding establishment of chemical equilibrium. Note that as the highly compressed hadronic system evolves in time, its entropy must increase or remain constant — typically it remains constant since further entropy production is difficult as model considerations show. More conveniently, one looks at the specific entropy per baryon $S/B$ which must nearly remain constant even in presence of considerable particle evaporation — such emission processes are likely to reduce the baryon and entropy content of the fireball at comparable rate. For physical systems such as is a gas of interacting hadrons, a time evolution at fixed $S/B$ is not compatible with a fixed value of $\lambda_q$ and, moreover, one should expect, in models in which particle freeze-out occurs subject to their respective interaction cross sections, different freeze out values of $\lambda_q$. Thus the final state particle abundances for a given collision system should be described by values of $\lambda_q$ which vary from particle to particle.

The exception to these observations corresponds to a temporal evolution under equations of state in which aside of the dimensionless variables $\lambda_q$, $\lambda_s$ no dimensional scale occurs aside of $T$ — since the specific entropy $S/B$ is a dimensionless quantity it can in this case be only a function of $\lambda_q$, the $T$ dependence must cancel. The only known hadronic physical system which satisfies this condition is the QGP phase, and in order to give a unique value of $\lambda_q$ for all particles it must be hadronizing directly into decoupled hadronic particles, as we already
argued above. Thus presence of unique values of \( \lambda_i = \{ \lambda_s, \lambda_q \} \) (referred to as presence of chemical equilibrium) is proving to be a strong, even though indirect evidence for the occurrence of the deconfined state.

4. The QGP equations of state: We now develop a framework in which we can compute the expected relative abundance of strange particles emerging at the end of the nuclear collision in which QGP phase was formed. An essential ingredient in our model are QGP equations of state. Our assumption is that the gluons \( G \) and light quarks \( u, d, \bar{u}, \bar{d} \) are considered in full equilibrium (thermal and abundance) while we shall allow for suppression of the occupancy of \( s, \bar{s} \) quarks by the factor \( \gamma_s \). We will allow for ‘thermal’ masses of all these particles according to the relation:

\[
m^2_i = m^2_0 + (cT)^2,
\]

where in principle we have \( c^2 \propto \alpha_s \), but given the current uncertainty regarding the value of the coefficient \( c \) we shall simply explore its consequence in the domain \( c \sim 2 \) arising for \( \alpha_s \sim 1 \) in the standard formulas [19]. We take \( m^0_q = 5 \approx 0 \) MeV, \( m^0_s = 160 \) MeV, and \( m^0_G = 0 \).

It has been found in perturbative thermal QCD [20] which has been confirmed by more sophisticated lattice gauge numerical calculations that another, often more significant effect of the interaction is the reduction of the available degrees of freedom. We implement the following effective counting of gluon and quark degrees of freedom:

\[
\begin{align*}
g_G &= 16 \rightarrow 16 \left(1 - \frac{15\alpha_s}{4\pi}\right), \\
g_{i-T} &= 6 \rightarrow 6 \left(1 - \frac{50\alpha_s}{21\pi}\right), \\
g_{i-B} &= 6 \rightarrow 6 \left(1 - \frac{2\alpha_s}{\pi}\right),
\end{align*}
\]

where \( i = u, d, s \) and for quarks two factors are needed: the factor \( g_{i-T} \) controls the expression when all chemical potentials vanish (the \( T^4 \) term in the partition function for massless quarks) while \( g_{i-B} \) is taken as coefficient of the additional terms which arise in presence of chemical potentials. Note that we treat all three light quark flavors on the same footing and that in principle these corrections were established only in the limit \( m_i \ll T \), and hence not in the limit here considered when the thermal mass exceeds the temperature. We favor for the QCD coupling \( \alpha_s = 0.6 \) but we have explored the dependence of our results on variations in \( \alpha_s \). Given these interaction effects, numerical integration of the Bose/Fermi distributions for quarks/gluons allow us to obtain any physical property of the QGP. We have first made many studies to assess how the different uncertainties in the parameters and initial state hypothesis affect our results, and we believe that our key findings here reported are not significantly affected.

5. Discussion of results: For a given initial state \( E/B \), Fig. 1 presents the resulting constraint between the values of initial temperature \( T \) and the quark fugacity \( \lambda_q \). The experimental results correspond to the prior analysis ([17, 18] for S–S, [9, 19] for S–W/Pb). We were stunned to see that for the expected set of the QGP parameters we have obtained as result an excellent agreement with the results determined in the prior phenomenological analysis of the particle spectra regarding \( T \) and from relative particle abundance regarding \( \lambda_q \).
in the initial state. We were thus encouraged to proceed to study the strangeness yield along the constant entropy trajectories which are shown for the S–W/Pb case in the $\mu_B$–$T$ plane in Fig. 2: the short-long dashed boundary arising for a given $E/B = 8.8$ GeV (assuming equal baryon and energy stopping, and also $\gamma_s = 0.8$), the evolution at $S/B = 40$ (initial $\lambda_q \simeq 1.64$, top solid line), 50 ($\lambda_q \simeq 1.48$, the likely experimental case, middle solid line), 60 ($\lambda_q \simeq 1.38$, bottom solid line, corresponding nearly to the S–S collisions) is as expected from prior qualitative considerations, along straight lines. They end near 150 MeV, corresponding to the hadronization picture developed previously [11].

We recall that the entropy content defines the hadronization multiplicity. To a good approximation one can divide the total entropy by the entropy per particle in a relativistic ($m \leq T$) Boltzmann gas, $S/N = 4$, in order to assess the number of particles emerging in the hadronization process. Two corrections are to be noted:

i) a heavy nonrelativistic hadronic resonance has a greater per particle entropy content ($S/N = 2.5 + m/T + \ldots$),

ii) resonance disintegration increases the number of final state particles.

These two effects largely cancel for hadronization at relatively small temperature and the relationship between the specific entropy and specific multiplicity remains approximately:

$$\frac{S}{B} \simeq 4 \frac{h}{B},$$

(3)
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(As usual, \( h \) denotes the total hadronic multiplicity). Thus the lines of specific entropy which correspond to fixed \( \lambda_q \) are leading to some determined specific multiplicity in the final state. Note that the higher \( \lambda_q \) (greater baryon content) corresponds to a smaller specific entropy content and thus a smaller specific multiplicity. This implies at equal initial specific energy that the initial observable temperature rises with \( \lambda_q \). Hence as we move from S–S/Ag to S–W/Pb and finally to Pb–Pb, the specific multiplicity, and hence \( D_Q \equiv (h^+ - h^-)/(h^+ + h^-) \) decreases (as usual \( h^i \) is the hadron multiplicity of charge \( i \)). This systematic behavior may appear surprising, since it implies that the hotter state has a smaller final per baryon particle multiplicity, but is completely logical and consistent considering that the initial condition is determined by specific energy which hardly changes, hence a baryon density increase implies energy density increase. This justifies the higher initial temperature which transfers a greater fraction of initial energy into transverse flow, resulting in smaller specific multiplicity. We note that in the S–Ag system \( \lambda_q \) is near the value 1.38 for S–S and certainly smaller than \( \lambda_q = 1.48 \) for S–W/Pb and hence it is at most 5 units of specific entropy different from the S–S system. This difference translates into a fraction (about 0.4) difference in \( h^- \) yield per participant (reduction for S–Ag as compared to S–S) and is thus hardly observable: this fact was noted already in the recent NA35 reports: the yield of \( h^-/B = 1.8 \pm 0.2 \) seen by NA35 [1], though 50% above the yield of HG [7] is the same for the S–S and S–Ag systems. However the value given is lower than the expectations based on the specific entropy content of the QGP phase discussed here, considering that \( S/B \simeq 60 \). There are a number of possible explanations of this discrepancy: the baryon content is overestimated due to substantial and asymmetric contamination by \( K^\pm \) of the central region; the formation of QGP in only a fraction of events or only in a fraction of the volume; etc..

This decrease in the specific entropy as we go to baryon richer systems is at the origin of the here important QGP result: we find that the ratio of strange quark abundance to the entropy content in the QGP phase is not \( T \) dependent, and moreover it depends little on the initial state, since the \( \lambda_q \) (= Const. at fixed specific entropy) dependence is much slower than the canceling \( T \)-dependence. We obtain:

\[
s_h \equiv \frac{\langle s + \bar{s}\rangle_{\text{QGP}}}{(S/4)_{\text{QGP}}} = 0.25 \cdot (1 + 0.47(\alpha_s - 0.6)) \cdot \gamma_s ,
\]

in the entire parameter region of interest to us.

We now relate the theoretical value \( s_h \) to experimental observables. By virtue of Eq.(3) the entropy content is related to the hadronic multiplicity generated in the interaction, and a suitable measure of the hadronic multiplicity is the sum of baryon content with the negative hadron multiplicity \( h^- \) multiplied by the factor 3 to allow for the positive and neutral particles:

\[
\frac{S}{4} \simeq h \simeq 3h^- + B .
\]

Aside of the primordial strangeness there is some contribution arising from gluon fragmentation [3]. Considering that each gluon hadronizes with 15% probability [4] into a \( ss \) pair we have:

\[
\langle s + \bar{s}\rangle = \langle s + \bar{s}\rangle_{\text{QGP}} \left( 1 + 0.3 \frac{(G)_{\text{QGP}}}{\langle s + \bar{s}\rangle_{\text{QGP}}} \right) .
\]
Since $\langle G \rangle_{\text{QGP}} / \langle s + \bar{s} \rangle_{\text{QGP}} = 0.52$ is akin to Eq. (4) practically independent of the statistical variables, gluon fragmentation increases the observable final state strangeness yield by $15\%$. This ‘experimental’ yield is approximately: $\langle s + \bar{s} \rangle \simeq 2K^0_s + K^+ + K^- + 2(\Lambda + \Sigma^0 + \Xi + \Xi^0)$.

We note that one can replace $K^+ + K^-$ by $2K^0_s$ as long as the isospin asymmetry is not too large. The factor 2 in front of the neutral hyperon and anti-hyperon abundance accounts qualitatively for the usually unobserved charged hyperons $\Sigma^\pm$, which account for about factor 1.6 and the higher strange particle resonances which remain unobserved, including hidden strangeness (e.g. $\eta, \eta'$); this ‘factor two’ rule has proven itself in A–A collisions$^\text{[10]}$. In view of this discussion we thus define:

$$s_{\text{h}}^{\text{exp}} = \frac{\langle s + \bar{s} \rangle}{3h^- + B}$$  \hspace{1cm} (7)

and proceed with our discussion as if $s_{\text{h}}^{\text{exp}} \simeq s_{\text{h}}$, see Eq. (4). We note here that the advantage of considering the total strangeness rather than the enhancement of some individual fraction, as is often common when discussing enhancements of kaons or hyperons, is that the sum of strangeness abundance, whatever the model used, is independent of the chemical composition of the source (i.e. of $\lambda_q$) and hence one can easier compare different collision systems, in which one expects different values of $\lambda_q$.

First, we recall the ‘background’ value found in N–N reactions (judiciously weighted p–n, p–p reactions) which follows from the results given in Table 11 of Ref. [16]. Using here Eq. (7), but omitting the $B$ term in the denominator we find $s_{\text{NN}} = 0.092 \pm 0.014$. It is unwise to interpret this result in terms of a thermal model, i.e. to use Eq. (4), but it is important for us to obtain a point to compare with our A–A results — we so find $\gamma_{s_{\text{NN}}} = 0.37 \pm 0.05$ — to be compared to the value $\gamma_{s_{\text{NN}}} = 0.22 \pm 0.05$ which was obtained using only strange particle yields $^\text{[17]}$. Both these results are shown in Fig. 3 where we show $\gamma_s$ for the different reaction systems as function of $\log_{10}(1.5A_p^{2/3}A_t^{1/3})$, projectile (p) and target (t), which is an accepted measure of the relative size of the interaction volume. The smaller N–N value is marked ‘s’ to stress that the hadronic multiplicity was not directly involved in its determination.

M. Gaździcki $^\text{[6]}$ has given detailed information regarding the multiplicity of strange particles found in S–S/Ag collisions — in portions of his analysis he used knowledge of the phase space distributions in one system to extrapolate the results of the other. We therefore combine the two sets of data to obtain $s_{\text{S}} = 0.14$, which may be interpreted in terms of Eq. (4) to imply: $\gamma_s = 0.56 \pm 0.07 \pm 0.15$. The second error is our estimate of the systematic error made in equating the here defined experimental value of $s_{\text{h}}$, Eq. (7), with the theoretically well defined quantity $s_{\text{h}}$, Eq. (4). The cross in Fig. 3 indicates this result, alongside with $\gamma_{s_{\text{SW}}} = 0.8 \pm 0.1$ which follows from the analysis $^\text{[10]}$ of S–W strange anti-baryon data $^\text{[3, 4]}$. The increasing trend of the strangeness occupancy as the size and presumably life-span of the system increases (see Fig. 3) and thus its qualitative agreement with the expected behavior for a system which is near but not at strangeness abundance equilibrium, is very encouraging — we note that as compared to the N–N results we have considerable rise in strangeness yield and that a total yield increase by a factor 2.5–5 is anticipated comparing Pb–Pb with N–N.

$^1$This prescription gives about 10% less strange particle yield compared to another procedure developed in $^\text{[21]}$. 

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Figure 3: The strangeness occupancy factor $\gamma_s$ as function of $\log_{10}(1.5A_p^{2/3}A_t^{1/3})$, (projectile (p) and target (t) nuclei). The cross is as determined here for the collision systems S–S/Ag. We also show prior results for S–W/Pb [10] and the two points for the N–N system which does not qualify for the thermal description: the result marked ‘s’ is based on strange particle yields [17], the other follows from the (non applicable) prescription here presented.

We note that in Ref. [17] the value $\gamma^{SS}_s = 1.0 \pm 0.2$ was obtained in a global particle abundance fit. This is not inconsistent with the current finding, as in this earlier approach a constraint was imposed that the relative abundance of hadronizing strange mesons and baryons follows the relative hadronic gas yield. Since, we have realized [11] that such a constraint is forcing a hadronization at $T = 190–200$ MeV, which indeed was found in Ref. [17]. In this condition no attempt could be made to understand the enhanced (50%) yield of predicted negative hadron multiplicity as compared to experiment. It is this unexplained $h^-$ enhancement which is at the origin of the here reported smaller value of $\gamma^{SS}_s$.

Very recently, the experiment NA36 has released its final experimental abundance for strange particles observed in S–Pb reactions [7]. However, the data does not cover the total strange particle abundance, but only particles in the interval $0.6 < p_\perp < 1.6$ GeV/c are given, and a further cut is made for rapidity window, with $1.5 < y < 3$. So restricted in this phase space window strangeness yield per total negative hadron multiplicity leads to $s_h^{cut} \sim 0.02$. This result is not inconsistent with the expectation $s_h^{Pb} = 0.8 \cdot 0.25 = 0.20$, but considerable further effort regarding extrapolation to the unobserved phase space regions is needed in order to compare theory with experiment.

We have shown in this paper that the hypothesis that QGP plasma state is formed is quantitatively in agreement with the observed strangeness enhancement in S–A collisions.
Our results may not be misconstrued to constitute a ‘proof’ of the formation of QGP, but should rather be seen as supportive evidence that the strangeness enhancement as reported is in agreement with the expectations based on our current knowledge about QGP form of hadronic matter. We have further found the interesting result that the observed primordial temperatures and the corresponding values of quark chemical fugacity are consistent with the notion of equal baryon number and energy stopping in S–A collisions, if QGP equations of state are assumed.

While previous strangeness enhancement analysis [21] has focused on the comparison of different strange particle yields with reference being the N–N collision system, we have established here a quantitative comparison with the expectations assuming the best current knowledge of the QGP state. In particular we are able to show that all experimental results including also the initial temperature seen at high $m_\perp$ are consistent with the QGP model in which we allow for a varying degree of strange phase space saturation, which is expected to approach unity as the size of the system increases. We find an increase in the value of $\gamma_s$ (see Fig. 3) as we move from S–S/Ag to S–W/Pb systems. Our finding of an increasing strangeness occupancy factor with increasing size of the interaction region proves that we can expect to reach full saturation in the case of the Pb–Pb collisions, for which we hence expect strange particle to all hadronic multiplicity ratio $s_{h \text{Pb}}^{\text{Pb}} = 0.25$, which implies that one in four particles produced will carry strangeness. It is also worthwhile to recall that we expect somewhat smaller particle multiplicity per participant and a significantly greater initial temperature in Pb–Pb collisions compared to S–W/Pb.

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References

[1] J. Rafelski, Phys. Rep. 88 (1982) 331.
[2] J. Rafelski and M. Danos, Phys. Lett. B192 (1987) 432.
[3] S. Abatzis et al., (WA85 collaboration) Phys. Lett. B259 (1991) 508;
S. Abatzis et al., (WA85 collaboration) Phys. Lett. B270 (1991) 123;
[4] D. Evans, for the WA85 collaboration, New results from WA85 on multi-strange hyperon production in 200 A GeV/c S–W interactions, presented at the “Quark Matter ’93” meeting, June 1993, to appear in the proceedings (Nucl. Phys. A).
[5] S. Abatzis et al., (WA85 collaboration) Phys. Lett. B316 (1993) 615.
[6] M. Gaździecki for the NA35 Collaboration, New Data on the Strangeness Enhancement in Central Nucleus-Nucleus Collisions at 200 A GeV, presented at the “Quark Matter ’93” meeting, June 1993, to appear in the proceedings (Nucl. Phys. A).
D. Röhrich for the NA35 Collaboration, it Hadron Production in S+Ag and S+Au
Strange Particle Abundance in QGP

Collisions at 200 GeV/Nucleon, presented at the “Quark Matter ’93” meeting, June 1993, to appear in the proceedings (Nucl. Phys. A).

[7] E. Anderson et al., NA36 Collaboration, Phys. Lett. B316 (1993) 603.

[8] J. Rafelski, Phys. Lett. B262 (1991) 333;
J. Rafelski, Nucl. Phys. A544 (1992) 279c, and references therein.

[9] J. Letessier, A. Tounsi, U. Heinz, J. Sollfrank and J. Rafelski, Phys. Rev. Lett. 70 (1993) 3530.

[10] J. Letessier, A. Tounsi, U. Heinz, J. Sollfrank and J. Rafelski, Strangeness Conservation in hot nuclear fireballs, Paris preprint LPTHE/92–27R submitted to Phys. Rev. D

[11] J. Letessier, J. Rafelski and A. Tounsi, Strange Particle Freeze-out, Paris preprint LPTHE/93–54, Phys. Lett. B, in press.

[12] J. Rafelski, On the trail of QGP: Strange Antibaryons in Nuclear Collisions, p.529 in: Particle Production in Highly Excited Matter, H. H. Gutbrod and J. Rafelski, NATO Physics series Vol. B 303, Plenum Press, New York 1993.

[13] R. Santo et al., for WA80 collaboration , Single Photon and neutral meson data from WA80 Preprint IKP-MS-93/0701, Univeristät Münster, Institut für Kernphysik.

[14] E. Schnedermann, J. Sollfrank and U. Heinz, Fireball Spectra, p.175, in: Particle Production in Highly Excited Matter, H. H. Gutbrod and J. Rafelski, NATO Physics series Vol. B 303, Plenum Press, New York 1993.

[15] I. Otterlund, Physics of Relativistic Nuclear Collisions, p.57, in: Particle Production in Highly Excited Matter, H. H. Gutbrod and J. Rafelski, NATO Physics series Vol. B 303, Plenum Press, New York 1993.
S.P. Sorensen et al. Nuclear Stopping Power, in: Proc. XXI Int. Symposium on Multiparticle Dynamics 1991, World Scientific, Singapore 1992.

[16] M. Gaździcki and O. Hansen, Nucl. Phys. A528 (1991) 754.

[17] J. Sollfrank, M. Gaździcki, U. Heinz, and J. Rafelski, Regensburg preprint TPR-93-14, Z. Physik C, in press.

[18] J. Rafelski and M. Danos, Strangeness Flow Differences in Nuclear Collisions at 15 and 200 GeV A, submitted to Phys. Rev.

[19] V. Klimov, Sov. J. Nucl. Phys. 33 (1981) 934;
H.A. Weldon, Phys. Rev. D26 (1982) 1394 and D26 (1982) 2789;
E. Braaten and R. Pisarski, Nucl. Phys. B337 (1990) 569 and 339 (1990) 310.

[20] S.A. Chin, Phys. Lett. B78 (1978) 552 and B119 (1982) 51.

[21] H. Białkowski, M. Gaździcki, W. Retyk and E. Skrzypczak, Z. Physik. C55 (1992) 491.