Chapter 31
Strangeness in Quark–Gluon Plasma – 1982

Johann Rafelski

Abstract It is argued that observation of the strange-particle abundance may lead to identification of the quark–gluon plasma and measurement of some of its properties. Approach to chemical equilibrium and competitive processes in the hadronic gas phase are discussed.

31.1 Overview

I would like to argue in this paper that the nature of the properties of quark–gluon plasma can be studied by observing the abundance of strange particles created in nuclear collisions, [1]. Unlike hadron–hadron collisions, we anticipate that in an important fraction of nucleus–nucleus collisions, each participating quark will scatter many times before joining in an asymptotic hadronic state. The associated simplification of the physics involved arises because the well-established methods of statistical physics can be used in such a case in order to connect the microscopic world with effects and properties visible to the experimentalist's eyes. Only the presumption of an approximate thermochemical equilibrium to be studied below in more detail, frees us from the dependence on details of quark wave functions.

As a consequence of the statistical equilibrium the available energy is equipartitioned among accessible degrees of freedom and, among other $s\bar{s}$ pairs. This means that there exists a domain in space in which, in a proper Lorentz frame, the energy of the longitudinal motion has been largely transformed to transverse degrees of freedom. The basic question concerns the internal structure of this hadronic fireball: instead of consisting of individual hadrons, it may be formed by quarks and gluons. In this new physical phase, these colour-charged particles are deconfined and can

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move freely over the volume of the fireball. It appears that the phase transition from the hadronic gas phase to the quark–gluon plasma is mainly controlled by the energy density of the fireball. Several estimates (See, e.g., [2]) lead to 0.6–1 GeV/fm^3 for the critical energy density, to be compared with around 0.25 GeV/fm^3 inside individual hadrons. Many theoretical questions about strong interactions will be settled once the parameters and nature of the phase transition have been determined.

Further development of this new field of research depends on the ability to observe plasma creation and its detailed physical properties. It is quite difficult to insert a thermometer and to measure baryon density at $T = 150$ MeV. We must either use only electromagnetically interacting particles [3] (photons, lepton pairs) in order to get them out of the plasma, or study the heavy quark flavour abundance, in particular strangeness, generated in the collision [1]. To obtain a better impression of what is meant, imagine that strange quarks are very abundant in the plasma (and indeed they are!). Then, for example, since the (*sss*)-state is bound and stable in the hot perturbative QCD vacuum, it would be the most abundant baryon to emerge from the plasma. I doubt that such an omegazation of nuclear matter could leave any doubts about the occurrence of the phase transition. But even the enhancement of the more accessible abundance of $\overline{\Lambda}$ may already be sufficient for our purposes.

I will now explain in more detail why the strange-particle abundance is so useful [1] for observing properties of the quark–gluon plasma. First we note that, at a given temperature, the quark–gluon plasma will contain an equal number of strange (s) and antistrange ($\overline{s}$) quarks, naturally assuming that the hadronic collision time is much too short to allow for light-flavour weak-interaction conversion to strangeness. Thus, assuming equilibrium in the quark plasma (see Sect. 31.2), we find the density of the strange quarks to be (two spins and three colours)

$$\frac{s}{v} = \frac{\overline{s}}{v} = 6\int \frac{\mathrm{d}^3 p}{(2\pi)^3}\frac{1}{\exp(\sqrt{p^2+m_s^2}/T)+1} \approx 3\frac{Tm_s^2}{\pi^2}K_2(m_s/T)\ , \tag{31.1}$$

neglecting for the time being the perturbative corrections. The mass m_s of the strange quarks in the perturbative vacuum is believed to be of the order of 180–300 MeV.[1] Since the phase space density of strangeness is not too high, the Boltzmann limit is used in Eq. (31.1). Similarly, there is a certain light antiquark density ($\overline{q}$ stands for either $\overline{u}$ or $\overline{d}$):

$$\frac{\overline{q}}{v} = 6\int \frac{\mathrm{d}^3 p}{(2\pi)^3}\frac{1}{\exp(|p|/T+\mu_q/T)+1} \approx \mathrm{e}^{-\mu_q/T}T^3\frac{6}{\pi^2}\ , \tag{31.2}$$

where the quark chemical potential is $\mu_q = \mu_\mathrm{B}/3$ and μ_B is the baryon chemical potential, we drop below the subscript 'B'. This exponent suppresses the $q\bar{q}$ pair production. It reflects the chemical equilibrium between q and $\bar{q}$ and the presence of a light quark density associated with the net baryon number.

[1]The 2014 reference value is $m_s(\mu = 2\,\mathrm{GeV}) = 95 \pm 5\,\mathrm{MeV}$.

Alternative, but physically equivalent, ways to understand these factors are the following two statements:

- $\bar{q}$ is Fermi-blocked, since in its production the partner q-quark has to go on top of a Fermi sphere ($T \to 0$ limit).
- $\bar{q}$ quarks are easily destroyed by the abundant q quarks in the plasma.

What we now intend to show is that there are often more $\bar{s}$ quarks than antiquarks of each light flavour. Indeed,

$$\frac{\bar{s}}{\bar{q}} = \frac{1}{2}\left(\frac{m_s}{T}\right)^2 K_2\left(\frac{m_s}{T}\right) \mathrm{e}^{\mu/3T} . \qquad (31.3)$$

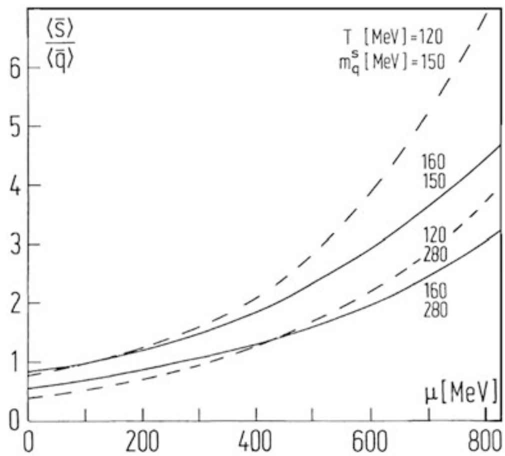
This ratio is shown in Fig. 31.1. Thus, we almost always have more $\bar{s}$ than $\bar{q}$ quarks and, in many cases of interest, $\bar{s}/\bar{q} \approx 5$. As $\mu \to 0$, there are about twice as many $\bar{u}$ or $\bar{d}$ quarks as there are $\bar{s}$ quarks at $T \gtrsim m_s$.

When the quark matter dissociates into hadrons, some of the numerous $\bar{s}$ quarks may, instead of being bound in a $q\bar{s}$ kaon, enter into a $(\bar{q}\,\bar{q}\,\bar{s})$ or $(\bar{q}\,\bar{s}\,\bar{s})$ antibaryon and, in particular, a $\overline{\Lambda}$, $\overline{\Sigma}$, or $\overline{\Xi}$. The probability for this process seems to be comparable to the similar one for the production of $\overline{\Lambda}$, $\overline{\Sigma}$, $\overline{\Xi}$, or $\overline{\Omega}$ by the quarks present in the plasma. What is particularly noteworthy about the $\bar{s}$-carrying antibaryons is that they can conventionally only be produced in direct pair production reactions. Up to high energies, this process is suppressed by energy-momentum conservation and phase space considerations. This leads me to argue that a study of the $\overline{\Lambda}$, $\overline{\Sigma}$, $\overline{\Xi}$, and $\overline{\Omega}$ in high energy nuclear collisions could shed light on the early stages of the nuclear collisions in which quark matter may be formed.

As is apparent from the above remark, the crucial aspects of the proposal to use strangeness as a tag of quark–gluon plasma involve:

- assumption of thermal *and* chemical equilibrium (see next section),

Fig. 31.1 Relative abundance of antistrange quarks $\bar{s}$ to light antiquark $\bar{q}$ as a function of μ for $T = 160\,\mathrm{MeV}$ (*solid lines*) and $T = 120$ (*dashed lines*), and strange quark mass $m_s = 150$, and 280 MeV, respectively



- comparison between results anticipated in both hadronic phases at given T and μ, the chemical potential to be determined by other considerations (see Sect. 31.3).

The theoretical techniques required for the description of the two quite different hadronic phases, the hadronic gas and the quark–gluon plasma, must allow for the formation of numerous hadronic resonances, which then dissolve at sufficiently high partial density in the state consisting of its constituents. At this point we must appreciate the importance of, and help provided by finite temperature. To obtain high particle density we may, instead of compressing the matter (which as it turns out is quite difficult), heat it up; many pions are generated easily, leading to the occurrence of the transition at moderate (even vanishing) baryon density [1].

31.2 Strangeness Production in the Quark–Gluon Plasma

In this section, we investigate the abundance of strangeness as a function of the lifetime and excitation of the plasma state [4]. This investigation was motivated by the observation that light quarks could not by themselves lead to chemical equilibrium of strange quarks [5]. After identifying the strangeness-producing mechanisms, we compute the relevant rates as a function of the energy density ('temperature') of the plasma state and compare them with those for light u and d quarks.

In lowest order in perturbative QCD, $s\bar{s}$ quark pairs can be created in collisions of two gluons (Fig. 31.2a–c) and by annihilation of light quark–antiquark pairs (Fig. 31.2d). The averaged total cross-sections for these processes were calculated by B. Combridge [6]. For fixed invariant mass-squared $s = (k_1 + k_2)^2$, where k_i are the four-momenta of the incoming particles, below $w(s) = (1 - 4M^2/s)^{1/2}$,

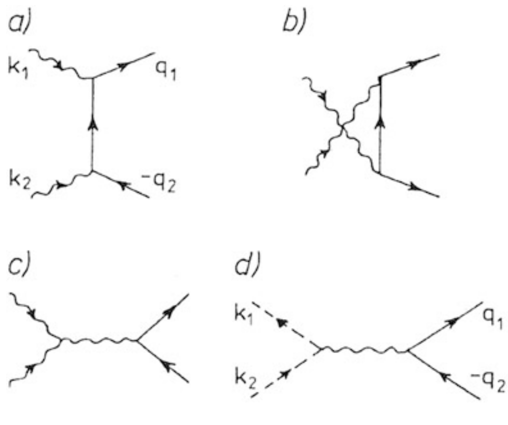


Fig. 31.2 Lowest order QCD diagrams for $s\bar{s}$ production: (**a–c**) $gg \rightarrow s\bar{s}$, and (**d**) $q\bar{q} \rightarrow s\bar{s}$

$$\overline{\sigma}_{gg\to s\bar{s}} = \frac{2\pi\alpha_s^2}{3s}\left[\left(1+\frac{4M^2}{s}+\frac{M^4}{s^2}\right)\tanh^{-1}w(s) - \left(\frac{7}{8}+\frac{31}{8}\frac{M^2}{s}\right)w(s)\right] . \tag{31.4a}$$

$$\overline{\sigma}_{q\bar{q}\to s\bar{s}} = \frac{8\pi\alpha_s^2}{27s}\left(1+\frac{2M^2}{s}\right)w(s) , \tag{31.4b}$$

For the mass of the strange quark, we explore the following cases:

- the value fitted within the MIT bag model: $M = 280$ MeV, and
- the value found in the study of quark currents: $M = 150$ MeV.

When discussing light quark production, we use $M = 15$ MeV. The effective QCD coupling constant $\alpha_s = g^2/4\pi$ is an average over space- and timelike domains of momentum transfers in the reactions shown in Fig. 31.2. We use (a) $\alpha_s = 2.2$, the value consistent with $M = 280$ MeV in the MIT bag model, and (b) the value $\alpha_s = 0.6$, expected at the involved momentum transfer, together with $M = 150$ MeV.

Given the averaged cross-sections, it is easy to calculate the rate of events per unit time, summed over all final and initial states[2]:

$$\frac{\mathrm{d}N}{\mathrm{d}t} = \frac{1}{2}\int \mathrm{d}^3x \int \frac{\mathrm{d}^3k_1}{(2\pi)^3|k_1|}\sum_i \rho_i(k_1,x)\int \frac{\mathrm{d}^3k_2}{(2\pi)^3|k_2|}\sum_i \rho_i(k_2,x) I k_1^{\mu} k_{2\,\mu}\overline{\sigma}(s). \tag{31.5}$$

The sum over initial states involves the discrete quantum numbers i (colour, spin, etc.) over which Eqs. (31.4a) and (31.4b) are averaged. The factor $k_1 \cdot k_2/|k_1||k_2|$ is the relative velocity for massless particles. We introduced a dummy integration $I \equiv \int_{4M^2}^{\infty} \mathrm{d}s\,\delta\big(s-(k_1+k_2)^2\big) = 1$ in order to facilitate the calculations. We now replace the phase space densities $\rho_I(k,x)$ by momentum distributions $f_g(k), f_q(k)$, and $f_{\bar{q}}(k)$ of gluons, quarks, and antiquarks that can still have a parametric x dependence, i.e., through a space dependence of temperature $T = T(x)$. The (invariant) rate per unit time and volume for the elementary processes shown in Fig. 31.2 is then

$$A = \frac{\mathrm{d}N}{\mathrm{d}t\mathrm{d}^3x} = \frac{1}{2}\int_{4M^2}^{\infty} s\mathrm{d}s\,\delta\big(s-(k_1+k_2)^2\big)\frac{\mathrm{d}^3k_1}{(2\pi)^3|k_1|}\frac{\mathrm{d}^3k_2}{(2\pi)^3|k_2|} \tag{31.6}$$
$$\left[\frac{(2\times 8)^2}{2}f_g(k_1)f_g(k_2)\overline{\sigma}_{gg\to s\bar{s}} + 2(2\times 3)^2 f_q(k_1)f_{\bar{q}}(k_2)\overline{\sigma}_{q\bar{q}\to s\bar{s}}\right] ,$$

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

[2] An additional factor 1/2 in the gluon production term is included in this printing: the wave function of two identical particles comprises the normalization factor $1/\sqrt{2}$ which when squared leads to $1/2$ in the thermal rate.

We furthermore assume that in the rest frame of the plasma, the distribution functions f depend only on the absolute value $|k| = k_0 \equiv k$ of the momentum. We then evaluate angular integrals in Eq. (31.6):

$$A = \frac{4}{\pi^4} \int_{4M^2}^{\infty} s\mathrm{d}s\, \overline{\sigma}_{gg\to s\bar{s}} \left[\int_0^{\infty} \mathrm{d}k_1 \int_0^{\infty} \mathrm{d}k_2 \Theta(4k_1k_2 - s) f_g(k_1) f_g(k_2) \right]$$
$$+ \frac{9}{4\pi^4} \int_{4M^2}^{\infty} s\mathrm{d}s\, \overline{\sigma}_{q\bar{q}\to s\bar{s}} \left[\int_0^{\infty} \mathrm{d}k_1 \int_0^{\infty} \mathrm{d}k_2 \Theta(4k_1k_2 - s) f_q(k_1) f_{\bar{q}}(k_2) \right], \tag{31.7}$$

where the step function Θ requires that $k_1k_2 \geq s/4 \geq M^2$. We now turn to the discussion of the momentum distribution and related questions. We note that the anticipated lifetime of the plasma created in nuclear collisions is of the order $6\,\mathrm{fm}/c = 2 \times 10^{-23}$ s. After this time, the high internal excitation will most likely have dissipated to below the energy density required for the global restoration of the perturbative QCD vacuum state (See also [7, 8]). The transition between the hadronic and the quark–gluon phase is expected at an energy density of approximately $1\,\mathrm{GeV/fm}^3$. Under these conditions, it is possible to estimate that each perturbative quantum (light quark, gluon) in the plasma state will rescatter several times during the lifetime of the plasma. Hence the momentum distribution functions $f(p)$ can be approximated by the statistical Bose (Fermi) distribution functions:

$$f_g(p) \approx \left(\mathrm{e}^{\beta\cdot p} - 1\right)^{-1} \qquad \text{(gluons)}, \tag{31.8}$$

$$f_{q/\bar{q}}(p) \approx \left[\mathrm{e}^{\beta\cdot p}\lambda^{(\pm)} + 1\right]^{-1} \qquad \text{(quarks/antiquarks)}, \tag{31.9}$$

where $\beta \cdot p = \beta_0 E_p - \boldsymbol{\beta}\cdot\mathbf{p}$, $E_p \to |\mathbf{p}|$ for massless particles, $(\beta \cdot \beta)^{-1/2} = T$ is the temperature-like parameter, and $\lambda^{(\pm)}$ is the baryon number (antibaryon number) fugacity. In the rest frame of the plasma, $\beta \cdot p = E_p/T \to |\mathbf{p}|/T$. The distributions [Eqs. (31.8) and (31.9)] can only be taken seriously for $|\mathbf{p}|$ not much larger than T; to populate the high-energy tail of the distributions, too many collisions are required, for which there may not be enough time during the lifetime of the plasma. While in each individual nuclear collision, the momentum distribution may vary, the ensemble of many collisions may lead to better statistical distributions.

Finally, let us discuss the values of the fugacities $\lambda^{(\pm)}$ in Eq. (31.9). As quarks are brought into the reaction by the colliding nuclei, baryon number conservation makes it possible to relate the baryon density ν to the fugacities by integrating Eq. (31.9) over all momenta:

$$\nu(T, \lambda^+, \lambda^-) = \frac{1}{3}\times 12 \int \frac{\mathrm{d}^3p}{(2\pi)^3} \left[\left(\mathrm{e}^{|\mathbf{p}|/T}\lambda^+ + 1\right)^{-1} - \left(\mathrm{e}^{|\mathbf{p}|/T}\lambda^- + 1\right)^{-1} \right]. \tag{31.10}$$

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

$$\lambda^+ = \frac{1}{\lambda^-} = \mathrm{e}^{-\mu_q/T} \,, \tag{31.11a}$$

$$\nu(T,\mu_q) = \frac{2}{3\pi^2}\left[\mu_q^3 + \mu_q(\pi T)^2\right] . \tag{31.11b}$$

As long as gluons dominate the $s\bar{s}$ formation in plasma state, conditions at the phase transition, such as abundance of q and $\bar{q}$, will not matter for the $s\bar{s}$ abundances at times comparable to the lifetime of the plasma. Hence, for the purpose of this study, we will use the value $\mu_q = 300$ MeV in order to estimate the quark densities at given temperature. We can now return to the evaluation of the rate integrals in Eq. (31.7).

In the glue part of the rate A, Eq. (31.7), the k_1, k_2 integral can be carried out exactly by expanding the Bose function Eq. (31.8) in a power series in $\exp(-k/T)$:

$$A_g = \frac{4}{\pi^4} T \int_{4M^2}^{\infty} \mathrm{d}s\, s^{3/2} \overline{\sigma}_{gg\to s\bar{s}}(s) \sum_{n,n'=1} (nn')^{-1/2} K_1\left(\frac{(nn's)^{1/2}}{T}\right) . \tag{31.12}$$

In the quark contribution, an expansion of the Fermi function is not possible and the integrals must be evaluated numerically. It is found that the gluon contribution of Eq. (31.12) dominates the rate A. For $T/M \gtrsim 1$, we find

$$A \approx A_g = \frac{7}{6\pi^2}\alpha_s^2 M T^3 \mathrm{e}^{-2M/T}\left(1 + \frac{51}{14}\frac{T}{M} + \cdots\right) . \tag{31.13}$$

Examples for A at $\alpha_s = 0.6$ and $M = 150$ MeV is shown in Fig. 31.3a. We note that in general the invariant production rate rises rapidly with T.

The abundance of $s\bar{s}$ pairs driven by A cannot grow forever. At some point the $s\bar{s}$ annihilation reaction will deplete the strange quark population. The $s\bar{s}$ pair annihilations may not only proceed via the two-gluon channel, but instead through γg final states [9]. The noteworthy feature of such a reaction is the production of relatively high energy γ's at the fixed energy of about 700–900 MeV ($T = 160$ MeV). These γ's will leave the plasma without further interactions. To some degree, this process is stimulated by coherent glue emission.

In any case, the $s\bar{s}$ annihilation loss term is proportional to the square of the density n_s of strange and antistrange quarks. With $n_s(\infty)$ being the saturation density at large times, the following differential equation determines n_s as a function of time [10]:

$$\frac{\mathrm{d}n_s}{\mathrm{d}t} \approx A\left\{1 - \left[\frac{n_s(t)}{n_g(\infty)}\right]^2\right\} . \tag{31.14}$$

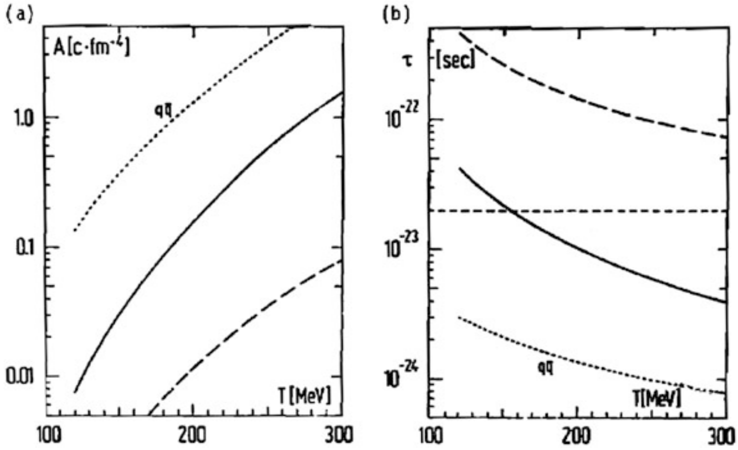


Fig. 31.3 (**a**) Rates A and (**b**) relaxation time constants τ, as a function of temperature T for $\alpha_s = 0.6$ and $M = 150\,\text{MeV}$. *Full lines* total process: $gg \to s\bar{s} + q\bar{q} \to s\bar{s}$. *Dashed lines*: $q\bar{q} \to s\bar{s}$. *Dotted lines*: $gg \to q\bar{q}$ where $M = 15$ MeV. *Note that on left rates A for glue based processes are shown too large by factor two*

The solution for initial value $n_s(t = 0) \approx 0$ is,

$$n_s(t) = n_s(\infty) \tanh \frac{t}{2\tau} \;\to\; n_s(\infty)\big(1 - 2e^{-t/\tau}\big) \,, \tag{31.15}$$

where[3] $\tau = n_s(\infty)/2A$. $n_s(t)$ is a monotonically increasing with temperature, saturating function, with asymptotic $t \to \infty$ behavior seen in Fig. 31.4b, controlled by the characteristic time constant τ. In a thermally equilibrated plasma, the asymptotic strangeness density $n_s(\infty)$ is that of a relativistic Fermi gas ($\lambda = 1$):

$$n_s(\infty) = \frac{2 \times 3}{2\pi^2} T M^2 \sum_{n=1}^{\infty} \frac{(-1)^{n-1}}{n} K_2(nM/T) \,, \tag{31.16}$$

provided the volume V is large.

For τ we find that the relaxation time[4]

$$\tau \approx \tau_g = \frac{9}{7}\left(\frac{\pi}{2}\right)^{1/2} \alpha_s^{-2} M^{1/2} T^{-3/2} \mathrm{e}^{M/T} \left(1 + \frac{99}{56}\frac{T}{M} + \cdots\right)^{-1} \tag{31.17}$$

is falling rapidly with increasing temperature as shown in Fig. 31.3b. While our results for strangeness production by light quarks agree in order of magnitude with

[3]The factor 2 in definition of τ had been inadvertently omitted.

[4]The following result has been always correct, however it combines two compensating, omitted factors 2 as noted above.

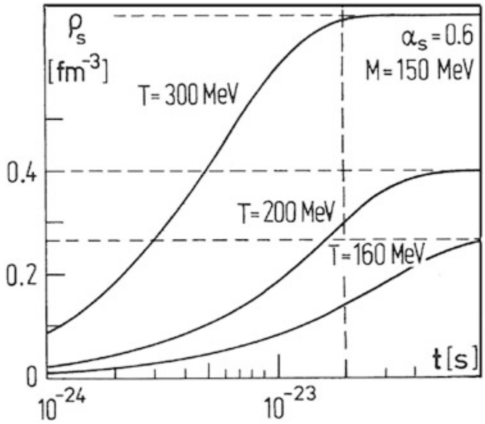


Fig. 31.4 Time-evolution of the strange quark density in the plasma for temperatures $T = 300, 200, 160\,\text{MeV}$ (*top to bottom*) for $m_s = 150$ MeV, $\alpha_s = 0.6$

those of Biró and Zimányi [5] (considering their choice of parameters), it is obvious here that gluonic strangeness production, which was not discussed by these authors, is the dominant process.

If we compare the time constant τ with the estimated lifetime of the plasma state, we find that the strangeness abundance will be chemically saturated for temperatures of 160 MeV and above, i.e., for an energy density above 1 GeV/fm^3. We note that τ is quite sensitive to the choice of the strange quark mass parameter and the coupling constant α_s and both must, however, be chosen consistently.

Also included in Fig. 31.3a,b are our results for gluon conversion into light quark–antiquark pairs. The shortness of τ for this process indicates that gluons and light quarks reach chemical equilibrium during the beginning stage of the plasma state, even if the quark/antiquark (i.e., baryon/meson) ratio was quite different in the prior hadronic compression phase.

The evolution of the density of strange quarks in Eq. (31.15) in the plasma state is shown in Fig. 31.4 for temperatures $T = 300, 200, 160\,\text{MeV}$. The saturation of the abundance is clearly visible for $T \geq 160\,\text{MeV}$. To obtain the measurable abundance of strange quarks, the corresponding values reached after the typical lifetime of the plasma state, 2×10^{-23} s, can be read off in Fig. 31.4 as a function of temperature. The strangeness abundance shows a pronounced threshold behaviour at $T \approx 120$–160 MeV.

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

31.3 Equilibrium Chemistry of Strange Particles in Hot Nuclear Matter

In order to establish the relevance of the strangeness signal for diagnosis of a possible formation of quark–gluon plasma, we must establish relevant particle rates originating from highly excited matter but consisting of individual hadrons—the hadronic gas phase, see [11] and Sect. 32.2. The main hypothesis which makes it possible to simplify the situation is to postulate the resonance dominance of hadron–hadron interactions (See, for example, [12])—in this case the hadronic gas phase is practically a superposition of an infinity of different hadronic gases and all information about the interaction is hidden in the mass spectrum $\tau(m^2, b)$ which describes the number of hadrons of baryon number b in a mass interval $\mathrm{d}m^2$ [13]. When considering strangeness-carrying particles, all we need to consider is the baryon chemical potential established predominantly by the non-strange particles.[5]

Here, we turn our interest directly to the rarest of all singly strange particles, and show in Fig. 31.5 the ratio $\langle n_{\overline{\Lambda}}\rangle / \langle n_{\Lambda}\rangle$. We notice an expected suppression of $\overline{\Lambda}$ due to the baryon chemical potential as well as strangeness chemistry. This ratio exhibits both a strong temperature dependence and a strong μ dependence. The remarkably small abundance of $\overline{\Lambda}$, e.g., $10^{-4}\Lambda$, under conditions likely to be reached in an experiment [$T \approx$ 120–180 MeV, $\mu \approx$ (4–6)T] is characteristic of the nuclear nature of the hot hadronic matter phase under consideration here. Our estimates for quark–gluon plasma based on flavour content are two to three orders of magnitude higher. We further note that $\mathrm{K}^+/\mathrm{K}^-$ abundance is a sensitive measure of the baryochemical potential, see Fig. 32.3.

In summary to this section, the relative abundance of strangeness-carrying antibaryons is greatly suppressed in the hadronic gas phase. Hence enhancements observed in nuclear collisions may be a useful indications of the reactions leading to the formation of the quark–gluon plasma. The study of multistrange hadrons is in progress.

31.4 Discussion

Only some selected aspects of the strangeness production in hot hadronic matter have been studied in detail. The results are quite encouraging and suggest interesting future perspectives. It was shown in Sect. 31.2 that strangeness abundance reaches chemical equilibrium in the plasma. The subsequent depletion of the strangeness during the plasma disintegration as well as its preferred hadronization channels have not yet been studied in detail. However, only if the plasma disintegration is an extremely slow process, lasting on the order of 10^{-22} s, can we anticipate a significant feedback on the high s abundance created at the maximum temperature reached in the collision. As shown in Fig. 31.3, the invariant rates drop quite rapidly with decreasing temperature, leading to a rapid increase in the equilibrium time constant τ; hence the strangeness abundance decouples from the equilibrium and remains a witness of the hot collision period.

While we cannot yet discuss in detail the abundance of multistrange antihadrons, which are influenced also by the possible ss, $\bar{s}\bar{s}$, sss, $\bar{s}\bar{s}\bar{s}$, and $s\bar{s}$ bound states

[5] To minimize duplication within this book we refer for the technical developments and the measurement of baryochemical potential μ to Sect. 32.2.

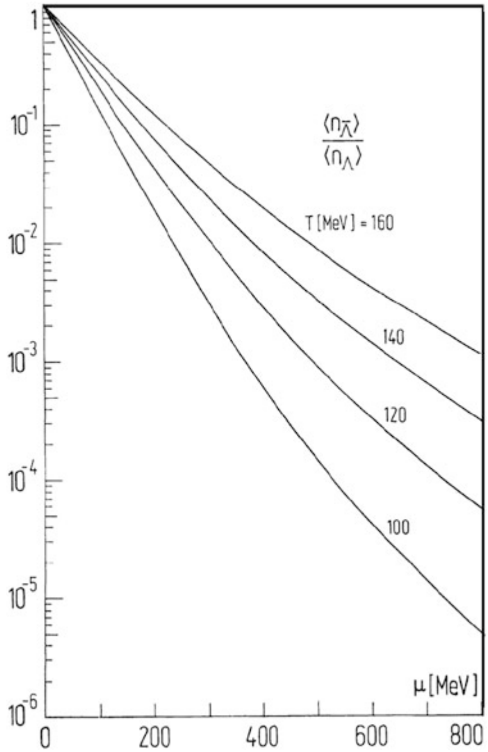


Fig. 31.5 The hadron gas ratio $\langle n_{\overline{\Lambda}} \rangle / \langle n_{\Lambda} \rangle$ as a function of μ for several temperatures $T = 100, 120, 140, 160\,\mathrm{MeV}$

in the plasma, it is apparent from the calculations performed in Sect. 31.3 that measurement of the production cross-section of the antistrange baryons could already be quite helpful in the observation of the phase transition. The high suppression of these degrees of freedom in the hadronic gas phase for obvious physical reasons is not maintained in the plasma phase, where $\bar{s}$ abundance is larger than $\bar{u}$, $\bar{d}$ abundance, as shown in Sect. 31.1. Measurement of the relative $\mathrm{K}^+/\mathrm{K}^-$ yield, while indicative for the value of the chemical potential, see Section 32.2, may carry less specific information about the plasma.

The K/π ratio may indeed also contain relevant information. However, it will be much more difficult to decipher the message. The π abundance will originate from diverse sources needed to be understood for that purpose.

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

It is important to appreciate that the experiments discussed above would certainly be complementary to the measurement with the help of electromagnetically interacting probes, e.g., dileptons or direct photons. Strangeness-based measurements have the advantage that they are based on the observation of a strongly interacting particle ($s, \bar{s}$ quark) originating from the hot plasma phase; these are much more abundant than the electromagnetic particles.

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