TOPICAL REVIEW

What is so special about strangeness in hot matter?†

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Abstract. The production of strange particles in a hot medium as produced in collisions of heavy ions is considered one of the most important signals for the phase transition to a quark-gluon plasma. In the first part of this lecture, the theoretical description of strangeness production in hot matter is outlined for a gas of quarks and gluons and for a hadronic gas and its impact on the deconfinement phase transition. Then in the second part, constraints from the underlying chiral symmetry of Quantum Chromodynamics (QCD) are utilized to extract signals with strangeness for the chiral phase transition in hot matter.

1. Introduction

This review intends to introduce to non-experts the basic ideas about the production of strange particles in hot, strongly interacting matter. I focus on matter at high temperatures and low baryon density, i.e. setting the baryochemical potential to zero. These conditions are most likely realized in the central region of the collision of heavy ions at relativistic bombarding energies where strongly interacting QCD matter can be probed in the laboratory. QCD exhibits a phase transition at high temperatures, around \( T_c \approx 170 \text{ MeV} \), as seen numerically on lattice gauge calculations (see e.g. [1]). Hadrons, mesons and baryons, are composite particles of quarks (and gluons) are present at low temperatures. Quarks and gluons are confined within the hadrons. Above the critical temperature, quarks and gluons are (asymptotically) free and not bound to hadrons anymore, i.e. a quark-gluon plasma (QGP) of deconfined quarks and gluons is formed.

The phase transition from a quark-gluon plasma to hadronic matter has happened during the early universe, about \( 10^{-4} \) s after the big-bang. Then the nucleons were formed during the deconfinement transition. In terrestrial laboratories one explores this phase transition by bombarding heavy nuclei at high energies and hunts for signals of the formation of the quark-gluon plasma (for a recent overview about the physics of the quark-gluon plasma, see [2]).

Particles with strangeness have been considered to be a particular useful probe of the quark-gluon plasma. There is a series of meetings dedicated to the topic of strange

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quarks in matter whose proceedings give an excellent overview of this field of research, see [3, 4, 5, 6, 7, 8, 9, 10, 11].

The enhanced production of strange particles was predicted to be a signal for the formation of a plasma of quarks and gluons in heavy-ion collisions (for a review on the physics of strangeness production in heavy-ion collisions see [12]). Indeed, recently, indications for the formation of a new form of strongly interacting matter, the strongly interacting quark-gluon plasma, in ultrarelativistic heavy-ion collisions have accumulated and strengthened so that the discovery of the quark-gluon plasma has been put forward by several of the most influential theoreticians in the field (see [13] for a list of review articles and [14, 15, 16, 17, 18]).

I will not touch the physics of strange matter at finite baryon density, in particular not the physics of the cold, dense quark matter with strangeness, its phenomenon of color superconductivity and the physics of strangeness in astrophysics. Here, I refer the interested reader to the extensive review articles about strange matter [19], about cold quark matter [20, 21] and about strangeness in astrophysics [22].

The paper is organized as follows: first, I will discuss the production of strange particles in a quark-gluon plasma, then its production in a hadronic gas at finite temperature. The focus in that section will be about the confined (hadron gas) and deconfined (quark-gluon plasma) phase. Then, I will address the issue of strangeness production in terms of symmetries of QCD, i.e. the chirally broken phase (hadrons) and the chirally restored phase (quarks and gluons). Both descriptions should be mutually compatible with each other, as one knows from lattice gauge simulations that both transitions, the deconfinement and the chiral phase transition, happen at the same critical temperature. Each section closes with a short discussion of recent developments in the corresponding research fields.

2. Strangeness and the deconfinement phase transition

In the following sections, I discuss the production mechanisms for producing strange particles for two distinctly different pictures: First, in the deconfined state of free quarks and gluons, and second, for a free hadron gas. Corrections due to interactions and dynamical effects are shortly addressed at the end of the corresponding subsections.

2.1. Strangeness in a quark-gluon plasma

In 1982, Rafelski and Müller demonstrated, that the production of strange quarks will be enhanced in heavy-ion collisions, if a plasma state of quarks and gluons is formed [23]. The arguments were basically twofold. First, the production threshold for the associated production of strangeness via a pair of strange-antistrange quark pairs is considerably smaller than the one for hadrons. Second, the equilibration timescale for producing strange particles in a quark-gluon plasma is much smaller than the one for a hadronic gas, so that the produced strange particles are not suppressed by dynamical
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Figure 1. The Feynman diagrams in perturbative QCD for the production of strange
and anti-strange quarks in a quark-gluon plasma.

effects and the corresponding number density is close to the equilibrium value.

The first argument is fairly easy to see. Consider the energy needed to produce
strange particles for a gas of quarks and gluons in comparison to the one for a hadron
gas. The associated production of a strange-antistrange quark pair can proceed by the
fusion of two gluons or two (massless) light quarks,

\[ q + \bar{q} \leftrightarrow s + \bar{s} \quad (q = u, d) \quad g + g \leftrightarrow s + \bar{s} \]

so that only an mass excess of

\[ Q_{qgp} = 2m_s \approx 200 \text{ MeV} \]

is involved. On the other hand, hadronic strangeness production proceeds in free space
via NN→NΛK with a considerably larger mass difference of the incoming and outgoing
hadrons (the Q-value) of

\[ Q_{hg} = m_\Lambda + m_K - m_N \approx 670 \text{ MeV} \]

Hence, strangeness production should be considerably enhanced in a quark-gluon plasma
relative to that of a free hadron gas.

For the second argument, a more elaborate and detailed calculation has to be
considered (see [24, 25] for details). The Feynman diagrams for the production of a
strange-antistrange quark pair are shown in Fig. 1 to first order in perturbation theory.
Note, that there are three diagrams for the production process via gluons, where one
is due to the nonabelian character of the gluons in QCD. Calculation of the Feynman
diagrams gives for the cross section involving quarks:

\[
\sigma_{q\bar{q} \rightarrow s\bar{s}} = \frac{8\pi\alpha_s^2}{27s} \left( 1 + \frac{2m_s^2}{s} \right) \left( 1 - \frac{4m^2_s}{s} \right)^{1/2} = \frac{8\pi\alpha_s^2}{27s^2} (s + 2m^2_s) \chi
\]

with

\[ \chi = \sqrt{1 - \frac{4m^2_s}{s}}. \]

For \( \alpha_s = 1/2 \) and a typical energy scale of \( s = (3T)^2 \approx (0.6\text{GeV})^2 \) in a thermal bath of
massless particles, one finds a cross section of about 0.25 mb for \( m_s = 100 \text{ MeV} \). The
 gluonic production processes result in a cross section of

\[
\sigma_{gg} = \frac{\pi\alpha_s^2}{3s} \left[ \left( 1 + \frac{4m^2_s}{s} + \frac{m^4}{s^2} \right) \ln \left( \frac{1 + \chi}{1 - \chi} \right) - \left( \frac{7}{4} + \frac{31m^2_s}{4s} \right) \chi \right],
\]
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Figure 2. The cross section for strange-antistrange quark pair production via gluons (solid lines) and quarks (dashed lines) in perturbative QCD for a strange quark mass of $m_s = 100$ MeV; upper lines are for $\alpha_s = 0.5$, lower lines for $\alpha_s = 0.3$.

which turns out to be 0.6 mb for the same parameters. Hence, gluon fusion is the dominant process for strangeness production in a quark-gluon plasma. The cross section as a function of energy is plotted in fig. 2 for two different values of the strong coupling constant, $\alpha_s = 0.5$ (upper lines) and $\alpha_s = 0.3$ (lower lines). The cross section has a threshold at $\sqrt{s} = 2m_s$, rises drastically, reaches a maximum just above the threshold and falls down rapidly. The gluon production cross section for strange-antistrange quarks dominates over the quark production cross section only for larger energies well beyond the maximum. Note, that the values of the perturbative cross sections are quite small, in the range of 1 mb and below, compared to a typical value of 40 mb for proton-proton collisions.

For the equilibration timescale, one has to look at the rate per unit time and volume in a heat bath of given temperature. The cross section has to be averaged over the distribution functions of the incoming particles:

$$A = \frac{dN}{dt d^3x} = \frac{1}{2} \int_{4m_s^2}^{\infty} ds \cdot s \cdot \delta \left( s - (k_1 + k_2)^2 \right)^2$$

$$\times \int \frac{d^3k_1}{(2\pi)^3|k_1|} \int \frac{d^3k_2}{(2\pi)^3|k_2|} \left\{ \frac{1}{2} (2 \times 8)^2 f_g(k_1) f_g(k_2) \sigma_g(s) \right\}$$

$$+ 2 \times (2 \times 3)^2 f_q(k_1) f_{\overline{q}}(k_2) \sigma_q(s)$$

which corresponds to the analogue of the thermal average of the cross section $\langle \sigma \cdot v \rangle$ in the nonrelativistic case. Here, $f_g$ and $f_q$ are the thermal distribution functions for gluons and quarks, respectively. Note, that no chemical potential is taken into account, the calculation assumes that there is zero baryon density in the hot medium. Also, Pauli-
blocking effects are ignored. They turn out to be small, as the phase space density of the produced strange quarks is not sufficiently high as to Pauli-block the reactions. Now, the density of produced strange quarks in the hot medium as a function of time can be derived by using the following master equation:

\[
\frac{dn_s}{dt} = A \cdot \left\{ 1 - \left( \frac{n_s(t)}{n_s(t = \infty)} \right)^2 \right\}
\]

where the strange quark density at infinite time \(n_s(t = \infty)\) corresponds to the equilibrium density of strange quarks \(n_{eq}^s\). The time evolution of the strange quark density eq. (6) is governed by a gain term and a loss term. The former one is just given by the rate \(A\), while the loss term is proportional to the squared density of already produced strange quarks. The normalization is chosen in such a way, that \(n_s\) saturates for infinite times at the equilibrium value. The equation can be formally solved for a constant rate to give

\[
n_s(t) = n_{eq}^s \cdot \tanh \left( \frac{t}{\tau_{eq}} \right)
\]

where \(\tau_{eq}\) stands for the equilibration time scale as defined by

\[
\tau_{eq} = \frac{1}{n_{eq}^s \cdot A}.
\]

The time to reach a certain equilibrium fraction \(f\) of strange quark density relative to the equilibrium one is determined by

\[
t_f = \tau_{eq} \cdot \tanh^{-1} \left( \frac{n_s}{n_{eq}} \right)
\]

This dependence of the equilibration time has to be compared to the standard ones for an exponential behaviour of the number density (like in ordinary radioactive decay) of \(t_f = -\ln(1 - n/\bar{n}_{eq})\). For a fraction of \(f = 1/e\), one gets \(t_f = \tau_{eq}\) and to reach an equilibrium fraction of \(f = 0.99\) it takes \(t_f = 4.6\tau_{eq}\) for the exponential case. For our case of strange-antistrange quark production in hot matter, which is proportional to the density squared, the times to reach a certain equilibrium fraction will be shorter, i.e. \(t_f \approx 0.74\tau\) for a fraction of \(f = 1/e\) and \(t_f \approx 2.6\tau\) for \(f = 0.99\). It is interesting to note, that a similar physical system which obeys the characteristics discussed here is the production of \(^3\)He in the proton-proton cycle in our sun whose production rate depends also on the density squared (see pp. 340 in [26]).

To arrive at absolute numbers for the equilibration time scale, the thermally averaged rate \(A\) has to be calculated explicitly. For a typical temperature of about \(T = 200\) MeV, the equilibration time turns out to be \(\tau_{eq}^{qgp} \approx 10\) fm, if a quark-gluon plasma is formed [23]. This equilibration time is about the timescale for a relativistic heavy-ion collision from the initial collisions until final freeze-out. Hence, it seems questionable that the system has enough time to bring the production of strange quarks close to its equilibrium value.

Fig. 3 shows the time evolution of the strange quark density for different choices of the temperature in the plasma. As one sees, the time to reach the equilibrium
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The number of strange quarks per baryon number in a hot gas of quarks and gluons as a function of the time. The solid horizontal lines denote the value in equilibrium for the different chosen temperatures. Reprinted figure with permission from [23]. Copyright (1982) by the American Physical Society.

Values depend strongly on the temperature. For \( T = 300 \text{ MeV} \), a time of about 3 fm seems to be enough to fully equilibrate strangeness production in the plasma, while at a temperature of \( T = 160 \text{ MeV} \), something like 20 fm are needed.

One should keep in mind, that it is assumed from the beginning, that an equilibrated quark-gluon plasma is formed with a given temperature. Of course, the quark-gluon plasma needs some time to be formed, as well as the temperature will drop when the system expands. Those dynamical effects are not taken into account in this simple estimate.

Moreover, which is even more severe, it is known from lattice data at finite temperature, that perturbation theory fails and can not describe the equation of state as extracted from the lattice. There are strong nonperturbative effects even for temperatures which are up to 4 times larger than the critical temperature, which has been measured to be \( T_c \approx 170 \text{ MeV} \) in full QCD (see e.g. [1]). Fig. 4 depicts the various contributions in perturbative QCD to the pressure as a function of temperature for each order. While the second order calculation computes a pressure which is below the one for an ideal gas of particles, the next order has a different sign and brings the total pressure well above the ideal gas pressure. Another change of sign occurs for the fifth order calculation which even brings the pressure below the result for the second order contribution. This oscillating behaviour and the fact that higher order contributions are larger than the lower ones, demonstrates clearly, that perturbation theory can not be used to describe the plasma of quarks and gluons not even at the phase transition but also not up to \( 4T_c \)!

Recent progress in finite temperature field theory utilizes resummation techniques to tackle this problem with considerable more success than pure perturbation theory like...
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![Figure 4](image.jpg)

Figure 4. The contribution to the pressure at finite temperature for different orders of perturbation theory in QCD. One notes that the different orders oscillate indicating a break-down of perturbation theory even at temperatures well above $T_c$. Reprinted figure with permission from [27]. Copyright (2000) by the American Physical Society.

Two-loop hard-thermal loop resummation [28] [29], for a review see [30]. Nonperturbative effects have been also included to some extent for the calculation of the production of strange quarks, like the cut-off model for gluons (see Fig. 11 in [31]), massive gluons [32], and resummation by hard-thermal loops using different approximation schemes [33] [34] [35]. Massive gluons allow for a new diagram of producing strange quarks, simply by the decay of a massive gluon to a strange-antistrange quark pair. On the other hand, a finite mass for gluons will suppress the gluon fusion processes. An overall consistent picture of the nonperturbative effects in equilibrium has not emerged yet, but the approaches studied so far indicate that the equilibration time scale will stay above 10 fm when nonperturbative effects are taken into account (see also the discussion in [36]). This picture is bolstered by the study of equilibration time scales for quarks and gluons in [37], where it was found that equilibration in a quark-gluon plasma is feasible for heavy-ion collisions at the future LHC but probably not at RHIC energies of $\sqrt{s} = 200$ GeV/nucleon.

There have been parton cascade models developed, which describe the formation and the non-equilibrium expansion of the quarks and gluons formed in a relativistic heavy-ion collision. The enhanced production of strangeness has been studied using the VNI [38] and the HIJING model [39]. Production of particles at large transverse momenta, where methods of perturbative QCD are applicable, and impacts for signaling the quark-gluon plasma are reviewed in detail by Gyulassy [40] taking into account nonperturbative effects (jet quenching) from a strongly interacting quark-gluon plasma.

With the advent of the data from RHIC, at least two other new paradigms of
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particle production substantiated: the notion of the formation of a saturated state of gluons, the colour glass condensate [41], and of quark coalescence [15]. The latter one is particular interesting for strange particles as they are more sensitive to collective effects and the signals proposed by quark coalescence. The relation of strange particle spectra and the colour glass condensate has been partially explored via a generalized $m_t$ scaling behaviour in [42].

It is clear from the above short discussion that strange quark production in hot QCD is far from being settled and still is a very active field of research.

2.2. Strangeness in hadronic matter

In this subsection, I discuss the production of strange particles in the hadronic picture. The stable baryons in the vacuum under strong interactions are besides the nucleons (N) the hyperons $\Lambda$ and $\Sigma^{+,-}\bar{0}$ with one strange quark, the $\Xi^{0,-\bar{0}}$ with two strange quarks, and the $\Omega^{-}$ with three strange quarks. The mass of the baryons increases with the number of strange quarks. The stable mesons under strong interactions are the pions $\pi^{+,-\bar{0}}$ and the kaons $K^{+,-}\bar{0}$ with one anti-strange quark and its antiparticle states $K^{-}$ and $K^{0}$. Besides those stable particles, there are more than hundred resonances known with a mass below 2 GeV which will also appear in hot matter and will form resonance matter. For our discussion, the first resonant state of the nucleon, the $\Delta(1232)$ will be especially important in our following discussion.

The production of strange particles in a free gas of hadrons has been studied in Koch, Müller and Rafelski [43]. The production of multiply strange baryons (the $\Xi$ and the $\Omega^{-}$) and of antihyperons was found to be particularly strongly suppressed in a hadronic gas, as the equilibration timescales for their production was much larger than the typical collision time of a heavy–ion collision.

The basic channels for strange hadron production in the vacuum are:

\[
\begin{align*}
\pi + \pi & \rightarrow K + \bar{K} \quad (Q = 2m_K - 2m_\pi \approx 710 \text{ MeV}) \\
N + N & \rightarrow N + \Lambda + K \quad (Q = m_\Lambda + m_K - m_N \approx 670 \text{ MeV}) \\
\pi + N & \rightarrow K + \Lambda \quad (Q = m_\Lambda + m_K - m_N - m_\pi \approx 530 \text{ MeV})
\end{align*}
\]

where the latter one can be studied by a secondary beam of pions in the laboratory. The Q-values for these processes are $Q = 710 \text{ MeV}$, $670 \text{ MeV}$, and $530 \text{ MeV}$, respectively, so substantially higher than for a quark-gluon plasma, where $Q = 2m_s \approx 200 \text{ MeV}$. In the hot medium, as stated above, resonances will appear so that channels like the following are possible:

\[
\begin{align*}
N + \Delta & \rightarrow N + \Lambda + K \quad (Q = m_\Lambda + m_K - m_\Delta \approx 380 \text{ MeV}) \\
\pi + \Delta & \rightarrow K + \Lambda \quad (Q = m_\Lambda + m_K - m_\Delta - m_\pi \approx 240 \text{ MeV})
\end{align*}
\]

Now the Q-value is already comparable to the one for a quark-gluon plasma! Even more, the Q-values can become smaller or even negative as for

\[
\Delta + \Delta \rightarrow N + \Lambda + K \quad (Q = m_\Lambda + m_K + m_N - 2m_\Delta \approx 90 \text{ MeV})
\]
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\[
\pi + \rho \rightarrow K + \bar{K} \quad (Q = 2m_K - m_\pi - m_\rho \approx 80 \text{ MeV}) \quad (16)
\]
\[
\rho + N \rightarrow K + \Lambda \quad (Q = m_\Lambda + m_K - m_N - m_\rho \approx -100 \text{ MeV}) \quad (17)
\]

Of course, the abundances of resonances like the \( \rho \) and the \( \Delta \) are suppressed exponentially by their (higher) masses. But the very low Q-values achievable for reactions with resonances means that it is possible that a resonance hadron gas might have production rates and equilibration timescales which are close to the one for a quark-gluon plasma. For a more quantitative answer one has to rely on a detailed numerical computation which includes all known resonances and their cross sections in hot matter.

Starting point is the evolution equation for the change of each particle number

\[
\frac{dN_i}{d^4x} = \sum_{j,k} <\sigma v>_i n_j(T)n_k(T) - \sum_l <\sigma v>_i n_i(T)n_l(T).
\]

(18)

The equation is a master equation consisting of a production term and a loss term. Again, Pauli blocking effects are not taken into account, as they are small corrections. The change of particle number is proportional to the density of the incoming particles and the thermally averaged cross section

\[
<\sigma v> \propto \int d^3p_1 d^3p_2 f_1(p_1)f_2(p_2)\sigma_{12}v_{12}
\]

(19)

where \( f \) are the distribution functions of the particles, i.e. Fermi-Dirac or Bose-Einstein distribution functions for thermally equilibrated matter. The cross section for most channels is poorly known if known at all, like for \( K + \Xi \rightarrow \Omega + \pi \), so that one has to assume some universal cross sections for those cases. The equilibration timescale one finds for a temperature of \( T = 160 \text{ MeV} \) are shown in Fig. 5 and can be read off to be about

\[
\tau_{\text{eq}}^h \approx 10^{-22} \text{ s} \approx 30 \text{ fm}
\]

for kaons, which is not so far from the one for a quark-gluon plasma. But the timescale for the (anti)hyperons, especially the \( \bar{\Xi} \) and \( \bar{\Omega} \) at finite density (finite baryochemical potential), can be an order of magnitude longer! The reason is, that it is more difficult to produce multiple units of strangeness in hadronic processes than in a quark-gluon plasma, where the produced strange quarks just coalesce to form a multiply strange baryon at particle freeze-out.

Detailed follow-up calculations in transport models supported this picture: the number of produced kaons in relativistic heavy-ion collisions can be explained by a purely hadronic picture at AGS energies \([14]\) and at SPS energies \([15]\). Note, that the often studied ratio of kaons to pions suffers from the fact that their nonlinear behaviour as a function of bombarding energy stems from the pions not the kaons \([15]\). Still, the kaon slopes can not be explained in present transport models \([16]\).

On the other hand, the production rates for antihyperons seems to be underestimated by hadronic transport simulations. A number of possible solutions have been proposed to remedy this, like colour ropes \([17]\) or multiquark droplets \([18]\), which rely on the quark-gluon picture. More recently, it was demonstrated, that the hadronic transport codes misses an essential reaction for a proper description of antibaryon
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Figure 5. The number density of various strange hadrons for a temperature of $T = 160$ MeV and different baryochemical potentials as a function of time. The solid horizontal lines denote the value in equilibrium for the different hadrons. Reprinted from "Strangeness in relativistic heavy ion collisions" [43], Copyright (1986), with permission from Elsevier.

production: the annihilation process of two antibaryons going to several mesons [49, 50]. The back reaction, several mesons produce a pair of antibaryons, is not a binary reaction, i.e. it involves more than two incoming particles, which is not taken into account in present transport codes (for a most recent attempt to incorporate $2 \leftrightarrow 3$ consistently see [51]). Nevertheless, the impacts of reactions like

$$p + \bar{p} \leftrightarrow n\pi \quad p + \bar{\Lambda} \leftrightarrow K + n\pi$$

can be studied using master equations as introduced above. The updated hadronic master equation models [49, 50] find now, that the production rate of antibaryons will be also enhanced in a purely hadronic approach by rescattering of multiple mesons into baryon-antibaryon pairs. Equilibration times for (anti)hyperons to be in chemical equilibrium can be as short as 10 fm as shown in [52]. Still, the production rates of the antihyperons $\Xi$ and $\Omega$ can not be described in this scenario unless effects from a phase transition and a rapid rise of the cross section with the pion number density out of equilibrium are added to the production rates [53].

An additional enhancement for strangeness production comes from the fact, that the hadron masses will experience medium modifications. In hot matter, the standard picture is that hadron masses decrease as a function of temperature which will lower the Q-value of the strangeness production processes even more and increase the number
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of produced strange particles. It is well known that \( \Lambda \) hyperons feel an attractive potential in nuclei, thereby forming bound states of a \( \Lambda \) and a nucleus, a so called hypernucleus (see e.g. [19] and references therein). Also, the antikaons will be modified substantially which has been studied in a coupled channel calculation for dense, cold matter [51, 53, 56, 57, 58] as well as for hot matter [59, 60].

Strange particle ratios, for antikaons but also those involving antihyperons, can be substantially enhanced in a hot medium in particular close to a phase transition (see e.g. [61, 60, 62] for calculations of particle ratios including in-medium effects). I will address medium modifications of hadrons in the next section utilizing chiral symmetry of QCD to describe hadron masses at finite temperature and its restoration as a measure of the phase-transition to a quark–gluon plasma.

3. Strangeness and Chiral Symmetry

In the previous section, I have been looking at a free gas of quarks and gluons in comparison to a free gas of hadrons. In this section, I want to focus on a completely different approach to study signals of the quark-gluon plasma with strange particles.

As pointed out earlier, it is known from lattice QCD simulations, that there is a phase transition at a temperature of \( T \approx 170 \text{ MeV} \) and that there are highly nonperturbative effects even well above that critical temperature. Lacking a detailed understanding of those nonperturbative features of QCD in hot and dense matter, one has to fall back on more fundamental features of QCD which do not need to incorporate a detailed treatment of the interactions explicitly: symmetries.

For the phase transition of pure gluon matter at finite temperature, lattice QCD simulations demonstrate that the deconfinement phase transition and the chiral phase transition coincide (see e.g. [63, 64]). Note, that one can assign an order parameter for the deconfinement phase transition, the Polyakov loop, only for the pure gluonic part of QCD. Once quarks are included, only the chiral order parameter, the quark condensate, remains as an order parameter which describes the chiral phase transition. Hence, chiral symmetry plays a crucial role in describing the phase transition of QCD in hot and dense, strongly interacting matter. Let us be more specific now and take a look at the QCD Lagrangian for three flavour massless quarks (\( q=\text{u,d,s} \) quarks):

\[
\mathcal{L}_{\text{qcd}}^0 = \bar{q}i\gamma_\mu (\partial^\mu - igA^\mu) q - \frac{1}{2} \text{Tr} G_{\mu\nu} G^{\mu\nu} .
\]  

The Lagrangian is invariant under the vector and axial transformations

\[
q' = q + i\alpha^a \frac{\lambda^a}{2} q \quad q' = q + i\beta^a \frac{\lambda^a}{2} \gamma_5 q
\]

with the corresponding conserved currents

\[
V_\mu^a = \bar{q}\gamma_\mu \frac{\lambda^a}{2} q \quad A_\mu^a = \bar{q}\gamma_\mu \gamma_5 \frac{\lambda^a}{2} q
\].
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The Lagrangian exhibits a chiral SU(3)$_L \times$ SU(3)$_R$ symmetry. One can assign now left-handed and right-handed quarks

$$q_L = \frac{1}{2} (1 - \gamma_5) q \quad q_R = \frac{1}{2} (1 + \gamma_5) q$$

which transform separately and do not mix with each other. Now in reality, the quarks have a finite mass which adds the following terms to the QCD Lagrangian:

$$\Delta L_{\text{mass}} = -m_u \bar{u}u - m_d \bar{d}d - m_s \bar{s}s$$

breaking the chiral symmetry explicitly. The mass terms mix the left-handed and the right-handed quarks as

$$\bar{q}q = \bar{q}_L q_R + \bar{q}_R q_L$$

Present estimates for the current quark masses range from

$$m_u = 1.5 - 5 \text{ MeV} \quad m_d = 3 - 9 \text{ MeV} \quad m_u = 60 - 170 \text{ MeV}$$

as given by the Particle Data Group [65]. Note, that the up and down quark masses are tiny in comparison to the nucleon mass $m_{u,d} \ll m_N \approx 1 \text{ GeV}$. The strange quark is much heavier than the light quarks but still considerably lighter than the nucleon or the hyperons, $m_s < 1 \text{ GeV}$. The tiny masses of the light quarks are actually essential to give the pion a finite mass. Nevertheless, a reasonable and rather successful assumption is to describe QCD by chiral symmetry plus corrections from explicit symmetry breaking.

Now, the masses of the hadrons are obviously not generated by quark masses in our world. The major contribution to the hadron masses comes from nonvanishing vacuum expectation values, i.e. from spontaneous chiral symmetry breaking. The Gell-Mann–Oaks–Renner (GOR) relation combines the nonvanishing expectation value for quarks, the quark condensate, with the pion mass and the pion decay constant $f_\pi = 92 \text{ MeV}$:

$$m_\pi^2 f_\pi^2 = -\frac{1}{2} (m_u + m_d) < \bar{u}u + \bar{d}d >$$

The relation can be motivated heuristically: the right hand side stems from the mass term of the QCD Lagrangian, the left hand side looks like a mass term for the pion where the field it couples to has a vacuum expectation value of just the pion decay constant. The equation then connects the quark world to the hadron world where the lightest known hadron is the pion. The GOR relation can be used to estimate the value of the quark condensate which is

$$< 0 | \bar{q}q | 0 > = - (310 \text{ MeV})^3$$

assuming an average light quark mass of 5 MeV. Also the gluon fields have a nonvanishing vacuum expectation value, the gluon condensate, which according to QCD sum rules for charmonium states [66] amounts to

$$< 0 | \frac{e_8}{\pi} G_{\mu\nu} G^{\mu\nu} | 0 > \approx (330 \text{ MeV})^4$$
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Figure 6. The lattice data for the light and strange quark condensate as a function of $\beta$ (a measure of the temperature). The data is taken from [64].

At first glance, the quark and gluon condensate seems to be of equal magnitude. However, the quark condensate has to be multiplied with the current quark mass so that

$$<0|{\bar{\alpha}_s \pi} G^a_{\mu\nu} G^{a\mu\nu}|0> \approx 80 <0|{\bar{q}q}|0> m_{u,d} \approx 4 <0|{\bar{q}q}|0> m_s$$

and the gluon condensate turns out to be much larger than the corresponding contribution from the quark condensate. It is the gluon condensate which generates the trace anomaly of QCD, i.e. a nonvanishing vacuum expectation value for the trace of the energy-momentum tensor

$$\theta^\mu_\mu = \frac{\beta_{\text{QCD}}}{2g} G^a_{\mu\nu} G^{a\mu\nu} \sum_{i=u,d,s} m_i \cdot \bar{q}_i q_i$$

where $\beta_{\text{QCD}}$ is the QCD $\beta$ function and $m_i$ the current quark mass. The trace anomaly can be related to the vacuum bag pressure [66], which basically determines the hadron masses in the MIT bag model.

Now the quark and gluon condensates will change in a hadronic medium. In particular at finite temperatures and zero density, the quark condensate is an order parameter for the chiral phase transition (strictly only in the chiral limit, i.e. for vanishing current quark masses). The quark condensate will melt, its value will drop to zero at the chiral phase transition temperature $T_\chi$. Lattice gauge simulations demonstrated that the light quark condensate indeed decreases drastically at a temperature of $T_c \approx 170$ MeV which coincides with the deconfinement phase transition.
[1] Lattice data also indicates that the strange quark condensate gets smaller at $T_c$, although much less pronounced due to the larger current mass of the strange quark [64]. Fig. 6 depicts the light quark and strange quark condensate as a function of $\beta$ which is a measure of the underlying temperature. One sees that the drastic drop of the light quark condensate at large temperatures (larger values of $\beta$) is accompanied by a moderate drop of the strange quark condensate. Hence, also strange particles will be moderated by the chiral or equivalently the deconfinement phase transition. Therefore, strange hadrons can in principle be utilized as a signal for the chiral phase transition. The advantage of using strange hadrons rather than light nonstrange ones is that strange particles can carry information from the high-density state as strangeness number is conserved in strong interactions.

The main task is now to relate the behaviour of the quark condensates to physical properties of (strange) hadrons. For that task one has to rely on effective models which incorporate the symmetry constrains of QCD. In the following, I will explore the effects of chiral symmetry restoration for strange hadrons in a SU(3)$\times$SU(3) chiral model [67] at finite temperatures. The hadrons involved are the pseudoscalar mesons ($\pi, K, \eta, \eta'$) and their chiral partners the scalar mesons ($\sigma, \kappa, a_0, f_0$) both forming a nonet in flavour SU(3). All 18 mesons can be grouped in one complex matrix

$$M = \Sigma + i \Pi = \sum_{a=0,8} \lambda_a (\sigma_a + i \pi_a)$$

(32)

where $\lambda_a$ denote the Gell-Mann matrices. One can form the following chiral invariants:

$$\text{Tr} M^\dagger M \quad \longrightarrow \quad O(18) \quad \text{(norm of vector)}$$

(33)

$$\text{Tr} M^\dagger MM^\dagger M \quad \longrightarrow \quad U(3) \times U(3) \quad (M \rightarrow UMU^{-1})$$

(34)

$$\det M + \det M^\dagger \quad \longrightarrow \quad SU(3) \times SU(3)$$

(35)

where the right side denotes the corresponding symmetry of the term. The last term breaks the $U_A(1)$ symmetry as $\text{det} \exp (i\lambda_0) = \exp (i \text{Tr} \lambda_0) \neq 1$. The model exhibits two order parameters, the expectation value of the $\sigma$ and the $f_0$ field ($\zeta$) which can be associated with the light quark and the strange quark condensate, respectively. In principle, there are three Gell-Mann matrices which are diagonal and can be associated with nonvanishing expectation values and order parameters, $\lambda_0, \lambda_3$ and $\lambda_8$, but $\lambda_3$ only breaks isospin symmetry and can be ignored in the following. The effective Lagrangian reads then

$$\mathcal{L} = \frac{1}{2} \text{Tr} \partial_\mu M^\dagger \partial^\mu M + \frac{1}{2} \mu^2 \text{Tr} M^\dagger M - \lambda \cdot \text{Tr} (M^\dagger MM^\dagger M)$$

$$- \lambda' \cdot (\text{Tr} M^\dagger M)^2 + c \cdot (\det M + \det M^\dagger)$$

(36)

which was used by Pisarski and Wilczek to study the order of the chiral phase transition in QCD [68]. The parameters of the model are $\mu^2, \lambda, \lambda'$, and $c$. They are determined by a fit to $m_\pi, m_K, m_\eta^2 + m_\eta'^2$, and $m_\sigma$. Explicit breaking of chiral symmetry is introduced in the model via

$$\mathcal{L}_{\text{esb}} = \epsilon \cdot \sigma + \epsilon' \cdot \zeta$$

(37)
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\[ \partial^\mu A_\mu^a = f_a m_\pi^2 a^a \]  

so that \( \langle \sigma_0 \rangle = f_\pi = 92.4 \) MeV and \( \langle \zeta_0 \rangle = \sqrt{2} f_K - f_\pi / \sqrt{2} = 94.5 \) MeV, where \( f_\pi \) and \( f_K \) are the pion and kaon decay constants, respectively.

The mass splittings in SU(3) are characterized in Fig. 7. In the chiral limit, i.e. for vanishing quark masses, there are nine Goldstone bosons and all nine pseudoscalar mesons of the nonet are massless. The physical \( \eta \) and \( \eta' \) are combinations of the singlet and octet state which mix in such a way that they are equal to the nonstrange \( \eta_{ns} \) and hidden strange \( \eta_s \) (this is the case of the so called ideal mixing). For a small but finite quark mass, there are nine pseudo–Goldstone bosons which have a finite mass. For the case that the coupling constants \( \lambda' \) and \( c \) are zero, all these pseudo–Goldstone bosons are degenerate, even for different light and strange quark masses. Letting \( \lambda' \neq 0 \) breaks the \( O(18) \) symmetry and the masses split according to their quark content. The pion and the \( \eta_{ns} \) are degenerate in mass, the kaon is heavier and the \( \eta_s \) as a pure \( \bar{s}s \)–state is the heaviest one. In the last column to the right, effects from the braking of the \( U_A(1) \) anomaly are switched on (\( c \neq 0 \)). Now the \( \eta' \) is a mixture of the \( \eta_{ns} \) and \( \eta_s \) and is the heaviest of all pseudoscalar mesons. The \( \eta \) mass is slightly larger than the kaon mass. What happened during the last stage is actually, that the degeneracy of the pion and the \( \eta_{ns} \) is lifted by the presence of the \( U_A(1) \) anomaly and the mass of the \( \eta_{ns} \) is shifted even above the mass of the kaon and \( \eta_s \). As the two \( \eta \) states mix with each other strongly, a level crossing occurs in reality so that there is a continuous line in mass shift between the \( \eta_s \) and \( \eta' \) and between the \( \eta_{ns} \) and the physical \( \eta \). Hence, the \( \eta \) and not the \( \eta' \) is expected to have a large (hidden) strangeness content!

The considerations above are made for zero temperature but give a quite accurate picture of what is going to happen at high temperatures when chiral symmetry is restored. If \( SU(2) \times SU(2) \) symmetry is restored, the chiral partners \( \pi–\sigma \) and \( \eta_{hs}–a_0 \) are degenerate in mass separately:

\[ m_\pi = m_\sigma < m_{\eta_{ns}} = m_{a_0} \]  

Figure 7. The mass splittings of the pseudoscalar mesons in the chiral SU(3) models.
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Figure 8. Meson masses as a function of temperature in an effective SU(3) chiral model. Note, that the mass for $\pi$ and $\sigma$ as well as for $a_0$ and $\eta$ are degenerate at $T_c \approx 150$ MeV and approach each other at larger temperatures (see [72]).

If in addition, the $U_A(1)$ anomaly is restored, these mass doublets will be equal in mass:

$$m_\pi = m_\sigma = m_{\eta_{ns}} = m_{a_0}$$  \hspace{1cm} (40)

If there is at least a partial restoration of the $U_A(1)$ symmetry at the critical temperature then

$$m_\pi = m_\sigma \approx m_{\eta_{ns}} = m_{a_0}$$  \hspace{1cm} (41)

The interesting and strange twist to this effect is, that the mass splitting of the purely nonstrange meson doublets is governed by the strange quark condensate, i.e. $\delta m^2 \propto c \cdot \zeta$!

That hadronic masses change in a hot medium has been known for quite some time and confirmed by lattice gauge simulations (see e.g. [69]). The pion mass rises as a function of temperature, while the mass of the $a_0$ drops for temperatures below $T_c$. The masses of the pion and the $a_0$ as well as for the pion and the kaon are degenerate at high temperatures well above $T_c$. It seems that the chiral anomaly is partially restored at $T_c$ as $m_{a_0}$ is close to $m_\pi$ at $T_c$ (see [70, 71, 72] and references therein). It is interesting to note, that the states are below the two–quark threshold so that the quark-gluon plasma seems to support hadronic excitations [69]! Even more than that, there seems to be bound states well beyond phase transition in the quark-gluon plasma at temperatures of say $T = (1-3)T_c$ [73].

The pseudoscalar and scalar mesons masses can be studied in the SU(3) linear sigma model at finite temperatures [72, 74]. Thermal fluctuations of the meson fields change
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Figure 9. Meson masses as a function of temperature in an effective SU(3) chiral model. The masses for the scalar kaon $K^*$ and the pseudoscalar kaon $K$ approach each other at $T_c$ and above while the $\eta'$ becomes even lighter than the kaon (see [72]).

The order parameters $\sigma$ and $\zeta$ which will change the meson masses. The whole set of equations for the thermal fluctuations, the order parameters and the meson masses can be solved selfconsistently at a given temperature. So far, the temperature dependence of the chiral anomaly has been put in by hand, so that the $U_A(1)$ symmetry is partially restored at $T_c$ and smoothly interpolates between the limiting cases of fixed anomaly coefficient $c \neq 0$ and $c = 0$. The effects of the chiral anomaly on the meson masses have been also studied in the Nambu–Jona-Lasinio model by Kunihiro [70, 71] with qualitative similar results.

Fig. 8 shows the (pole) mass for the nonstrange mesons $\pi$, $\sigma$, $\eta$ and $a_0$ as a function of temperature. At small temperatures, the meson masses hardly change. Only for temperatures close to $T \approx 150$ MeV, the masses of the scalar mesons $\sigma$ and $a_0$ drop drastically. The pseudoscalar meson masses increases slightly with temperature so that their curves merge with their corresponding chiral partners at $T = 150$ MeV. At this temperature, chiral symmetry for the light quark sector is restored (besides small effects from the explicit symmetry breaking). Still, there is a substantial mass splitting between the chiral doublets of $\pi - \sigma$ and $\eta - a_0$ due to the chiral anomaly which stems from the (still) nonvanishing strange quark condensate! It is only for much larger energies that all these nonstrange mesons become degenerate in mass.

Fig. 9 shows the temperature dependence of the meson masses of the kaon, $\kappa$ ($K^*$)
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and \( \eta' \) in comparison to the one for the pion. Like before, the mass of the scalar meson, the \( \kappa \), drops drastically, while the one for the kaon slightly increase so that their curves approach each other at \( T > 150 \text{ MeV} \). The mass of the \( \eta' \) decreases with temperature, contrary to the mass of the \( \eta \), and even gets lower than the one for the kaon. Note, that for temperatures larger than the crossover temperature of \( T = 150 \text{ MeV} \), the remaining mass splittings originates from the strange quark condensate which melts less steeply at \( T_c \) than the light quark condensate. For large temperatures, \( T \gg T_c \), all meson masses will become degenerate.

The drastic change of the mass spectra at finite temperatures should have observable consequences for relativistic heavy-ion collisions. Indeed, as the masses of the scalar mesons \( \sigma \), \( a_0 \) and \( \kappa \) fall so drastically close to the transition temperature, their decay channels to pseudoscalar mesons, \( \sigma \rightarrow \pi^+ \pi^- \), \( a_0 \rightarrow \eta + \pi \) and \( \kappa \rightarrow K + \pi \), will be closed simply by phase space arguments! The 'precursor' phenomenon of the chiral phase transition has been pointed out in [75] for the spectral function of the \( \sigma \) meson, the analogue for the \( a_0 \) and \( \kappa \) was studied in [72]. While the measurement of \( \eta' \)'s proceeds via leptonic decays and is difficult to address experimentally, the decay products of the \( \sigma \) and \( \kappa \) are readily measurable for a single collision event. Note, that the number of combinations for the kaon plus pion mass spectra is less than that for two pions as the kaon to pion ratio is less than one (about 0.17 for central gold-gold collisions) making the strange spectra more feasible experimentally.

The invariant mass spectra of two pions was measured for heavy-ion collisions of two gold nuclei at a bombarding energy of 200AGeV and used to extract the \( \rho \) meson [76] and the one for kaons plus pions to get the \( K^*(892) \) vector meson [77, 78]. So far, the \( \sigma \) meson as well as the \( \kappa \) meson has not been observed in these spectra. It will be difficult to extract clear signs for the appearance of the \( \sigma \) meson or the \( \kappa \) meson, as the resonances are only narrow around \( T_c \) and the decay products can rescatter and wash out any narrow resonance structure. On the other hand, fast hadronization at \( T_c \) will help to preserve the resonance in the mass spectra of correlated pions and kaons. One might have to wait for more complete and statistically precise data becoming available to clarify this issue.

Nevertheless, it is clear that the masses of strange mesons will be substantially reduced at the chiral phase transition, eventually being close to the masses of nonstrange mesons. Hence, again the quark-gluon plasma as being a chirally restored phase of strongly interacting matter has the appealing feature that the production of strange hadrons will be enhanced. Even more than that, there will be flavour equilibration, i.e. nonstrange and strange particle yields are similar to each other, as all masses are nearly degenerate above \( T_c \).

The questions remains how the chiral phase transition will influence the mass spectra of other hadrons, in particular the masses of (strange) vector mesons and baryons (hyperons) and its antiparticles. This issue has been addressed in a chiral SU(3) model at finite temperature and density [62]. The calculated particle ratios change drastically when medium modified masses are incorporated. The mass shifts are so strong that
only temperatures and densities below the chiral phase transition are compatible with
the experimentally measured particle ratios. In addition, a $\chi^2$ fit to the particle ratio
data demonstrates, that the freeze-out of particles in relativistic heavy-ion collisions in
equilibrium should then just happen at a temperature of a few MeV below the chiral
phase transition temperature. As the mass changes so drastically around $T_c$ it is unlikely
that the system can adjust in the short timescale of the collision and the particles have
to emerge out of equilibrium from the chirally restored phase.

4. Summary

In these lecture notes, the production of strange particles in a hot strongly interacting
medium have been described. Starting motivation for these investigations was the idea
that strange particles are more abundantly produced in hot matter of quarks and gluons,
the quark-gluon plasma. The topic has been tackled from quite different points of view:
first from the fact the quark-gluon plasma constitutes a state of deconfined matter,
second from the observations from lattice QCD that the phase transition at $T_c \approx 160$
MeV coincides with chiral symmetry restoration, i.e. the quark-gluon plasma is the
chirally restored phase of QCD.

In deconfined matter of quarks and gluons, the production rates for producing
strange-antistrange quark pairs should be larger compared to that for strange-
antistrange hadron pairs, as the threshold is considerably smaller. In addition, the
equilibration timescale for strangeness production in quark-gluon plasma should be
substantially smaller than the timescale of a typical relativistic heavy-ion collisions.
Estimates based on an ideal gas of quarks and gluons support the picture. However,
lattice gauge simulations demonstrate that there are strong corrections from the
non-ideal behaviour of a quark-gluon plasma even well above the phase transition
temperature. First attempts to incorporate nonperturbative effects find that the
equilibration timescale is shifted upwards close to the dynamical scale of heavy-ion
collisions.

The hadronic gas, for comparison, has to overcome a much larger Q-value for the
associated production of strangeness. This, however, holds strictly only for binary
collisions in free space. In a hot hadron gas, scattering of secondary particles, resonances,
lower the Q-value drastically, so that light strange hadrons, as the kaons, can be
produced more easily. Antibaryons and in particular antihyperons can be produced
by multi-pion and kaon fusion processes. The arguments can be extended to heavier
hadrons with multiple units of strangeness, as the $\Xi$ or the $\Omega^-$, but seem to fail to
describe the RHIC data if not additional effects from a phase transition are taken into
account.

The medium effects of (strange) hadrons have been studied in a chiral SU(3) model
for the pseudoscalar and scalar mesons. The hadron masses in vacuum are basically
§ This is at least correct for large temperatures and zero net baryon density, it need not be the case
for small temperatures and large densities, see the discussion in [79].
determined by vacuum expectation values of the quark and gluon condensates. For temperatures below $T_c$, the hadron masses hardly change as a function of temperature. However, close to $T_c$, the mass spectra dramatically changes. As chiral symmetry gets restored, the quark and gluon condensates melt. Then, the masses of the chiral partners of pseudoscalar and scalar mesons ($\pi - \sigma$, $\eta - a_0$, $K - \kappa$) have to become degenerate. There is some intrinsic mass splittings between some chiral partners due to the chiral $U_A(1)$ anomaly which is proportional to the strange quark condensate. The chiral anomaly diminishes as the strange quark condensate vanishes. The mass spectra at temperatures above $T_c$ approaches a flavour independent one, establishing flavour equilibration in the chirally restored phase. Calculation of a hot medium taking into account effects from chiral symmetry restoration finds that the hadron masses can not substantially change in order to get the experimentally observed particle ratios. This might indicate, that the particles have to freeze-out closely below the chiral phase transition.

It is interesting to see, how the two pictures of the phase transition of QCD seem to merge into an overall consistent picture. However, I think there needs to be a lot of work to be done and I hope this article will stimulate further research to finally settle the issue of strangeness production and the phase transition in hot and strongly interacting QCD matter.

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References

[1] F. Karsch, hep-lat/0401031 (2004).
[2] D. H. Rischke, Prog. Part. Nucl. Phys. 52, 197 (2004).
[3] Strange Quark Matter In Physics And Astrophysics, Nucl. Phys. B (Proc. Suppl.) 24B, 1 (1991), 20–24 May 1991, Aarhus, Denmark.
[4] International Symposium On Strangeness And Quark Matter, edited by G. Vassiliadis, A. D. Panagiotou, S. Kumar, and J. Madsen (World Scientific, River Edge, N.J., 1995), 1–5 September 1994, Chania, Crete, Greece.
[5] Strangeness In Hot Hadronic Matter: Strangeness '95, edited by J. Rafelski (American Institute of Physics, Woddbury, N.Y., 1995), 4–6 January 1995, Tucson, Arizona, USA.
[6] Strangeness In Hadronic Matter (Strangeness '96), Vol. 4 of Heavy Ion Physics, edited by T. Csörgő, P. Lévai, and J. Zimanyi (Akadémiai Kiadó, Budapest, 1996), 15–17 May 1996, Budapest, Hungary.
[7] International Symposium On Strangeness In Quark Matter, J. Phys. G 23, 1785 (1997), 14–18 April 1997, Santorini, Greece.
What is so special about strangeness in hot matter?

[8] *International Symposium On Strangeness In Quark Matter*, J. Phys. G 25, 143 (1999), 20–24 July 1998, Padua, Italy.
[9] *5th International Conference On Strangeness In Quark Matter*, J. Phys. G 27, 255 (2001), 20–25 July 2000, Berkeley, California, USA.
[10] *6th International Conference On Strange Quarks In Matter: A Flavourspace Odyssey (SQM2001)*, J. Phys. G 28, 1517 (2002), 24–29 September 2001, Frankfurt, Germany.
[11] *7th International Conference On Strangeness In Quark Matter (SQM2003)*, J. Phys. G 30, S1 (2004), 12–17 March 2003, Atlantic Beach, North Carolina, USA.
[12] C. Greiner, J. Phys. G 28, 1631 (2002).
[13] see http://quark.phy.bnl.gov/~mclerran/qgp/.
[14] M. Gyulassy and L. McMerran, nucl-th/0405013 (2004).
[15] B. Müller, nucl-th/0404015 (2004).
[16] E. V. Shuryak, hep-ph/0405066 (2004).
[17] H. Stöcker, nucl-th/0406018 (2004).
[18] X.-N. Wang, nucl-th/0405017 (2004).
[19] C. Greiner and J. Schaffner-Bielich, in *Heavy Elements and Related New Phenomena*, edited by W. Greiner and R. Gupta (World Scientific, Singapore, 1999), Vol. II, pp. 1074–1146, nucl-th/9801062.
[20] M. Alford, nucl-th/0312007 (2003).
[21] T. Schäfer, hep-ph/0304281 (2003).
[22] J. Rafelski and B. Müller, Phys. Rev. Lett. 48, 1066 (1982), erratum: ibid 56, 2334 (1986).
[23] R. D. Field, *Applications of Perturbative QCD* (Perseus Books Group, New York, 1989).
[24] B. L. Combridge, Nucl. Phys. B 151, 429 (1979).
[25] C. Rolfs and W. Rodney, *Cauldrons in the Cosmos* (The University of Chicago Press, Chicago and London, 1988).
[26] J. O. Andersen, E. Braaten, and M. Strickland, Phys. Rev. D 61, 014017 (2000).
[27] J. O. Andersen, E. Braaten, E. Petitgirard, and M. Strickland, Phys. Rev. D 66, 085016 (2002).
[28] J. O. Andersen, E. Petitgirard, and M. Strickland, hep-ph/0302069 (2003).
[29] J. O. Andersen and M. Strickland, hep-ph/0404164 (2004).
[30] W. Greiner and D. H. Rischke, Phys. Rep. 264, 183 (1996).
[31] T. S. Biró, P. Levai, and B. Müller, Phys. Rev. D 42, 3078 (1990).
[32] T. Altherr and D. Seibert, Phys. Lett. B 313, 149 (1993), erratum: ibid 316, 633 (1993).
[33] T. Altherr and D. Seibert, Phys. Rev. C 49, 1684 (1994).
[34] N. Bilic, J. Cleymans, I. Dadic, and D. Hislop, Phys. Rev. C 52, 401 (1995).
[35] J. Sollfrank and U. W. Heinz, in *Quark-Gluon Plasma II*, edited by R. C. Hwa (World Scientific, Singapore, 1995), p. 555.
[36] D. M. Elliott and D. H. Rischke, Nucl. Phys. A671, 583 (2000).
[37] K. Geiger and J. I. Kapusta, Phys. Rev. D 47, 4905 (1993).
[38] P. Csizmadia, P. Levai, S. E. Vance, T. S. Biró, M. Gyulassy, and J. Zimanyi, J. Phys. G 25, 321 (1999).
[39] M. Gyulassy, nucl-th/0403032 (2004).
[40] L. McLerran, hep-ph/0402137 (2004).
[41] J. Schaffner-Bielich, D. Kharzeev, L. D. McLerran, and R. Venugopalan, Nucl. Phys. A705, 494 (2002).
[42] P. Koch, B. Müller, and J. Rafelski, Phys. Rep. 142, 167 (1986).
[43] R. Mattiello, H. Sorge, H. Stöcker, and W. Greiner, Phys. Rev. Lett. 63, 1459 (1989).
[44] H. Weber, E. L. Bratkovskaya, W. Cassing, and H. Stöcker, Phys. Rev. C 67, 014904 (2003).
[45] E. L. Bratkovskaya, M. Bleicher, M. Reiter, S. Soff, H. Stöcker, M. van Leeuwen, S. Bass, and W. Cassing, Phys. Rev. C 69, 054907 (2004).
What is so special about strangeness in hot matter?

[47] H. Sorge, M. Berenguer, H. Stöcker, and W. Greiner, Phys. Lett. B 289, 6 (1992).
[48] K. Werner and J. Aichelin, Phys. Lett. B 308, 372 (1993).
[49] R. Rapp and E. V. Shuryak, Phys. Rev. Lett. 86, 2980 (2001).
[50] C. Greiner and S. Leupold, J. Phys. G 27, L95 (2001).
[51] C. Greiner, S. Juchem, and Z. Xu, hep-ph/0404022 (2004).
[52] J. Kapusta and I. Shovkovy, Phys. Rev. C 68, 014901 (2003).
[53] P. Braun-Munzinger, J. Stachel, and C. Wetterich, nucl-th/0311005 (2003).
[54] V. Koch, Phys. Lett. B 337, 7 (1994).
[55] T. Waas, N. Kaiser, and W. Weise, Phys. Lett. B 379, 34 (1996).
[56] M. Lutz, Phys. Lett. B 426, 12 (1998).
[57] A. Ramos and E. Oset, Nucl. Phys. A671, 481 (2000).
[58] L. Tolos, A. Ramos, and A. Polls, Phys. Rev. C65, 054907 (2002).
[59] J. Schaffner-Bielich, V. Koch, and M. Effenberger, Nucl. Phys. A669, 153 (2000).
[60] L. Tolos, A. Polls, A. Ramos, and J. Schaffner-Bielich, Phys. Rev. C 68, 024903 (2003).
[61] J. Schaffner, I. N. Mishustin, L. M. Satarov, H. Stöcker, and W. Greiner, Z. Phys. A 341, 47 (1991).
[62] D. Zschiesche, S. Schramm, J. Schaffner-Bielich, H. Stöcker, and W. Greiner, Phys. Lett. B 547, 7 (2002).
[63] J. B. Kogut, M. Stone, H. W. Wyld, S. Sheng, and D. K. Sinclair, Nucl. Phys. B 225, 326 (1983).
[64] J. B. Kogut, D. K. Sinclair, and K. C. Wang, Phys. Lett. B 263, 101 (1991).
[65] K. Hagiwara et al., Phys. Rev. D 66, 010001 (2002).
[66] M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, Nucl. Phys. B 147, 448 (1979).
[67] M. Gell-Mann and M. Levy, Nuovo Cimento 16, 705 (1960).
[68] R. D. Pisarski and F. Wilczek, Phys. Rev. D 29, 338 (1984).
[69] S. Gottlieb et al., Phys. Rev. D 55, 6852 (1997).
[70] T. Kunihiro, Phys. Lett. B 219, 363 (1989).
[71] T. Kunihiro, Nucl. Phys. B 351, 593 (1991).
[72] J. Schaffner-Bielich, Phys. Rev. Lett. 84, 3261 (2000).
[73] E. V. Shuryak and I. Zahed, hep-ph/0403127 (2004).
[74] J. T. Lenaghan, D. H. Rischke, and J. Schaffner-Bielich, Phys. Rev. D 62, 085008 (2000).
[75] T. Hatsuda and T. Kunihiro, Phys. Rev. Lett. 55, 158 (1985).
[76] J. Adams et al., Phys. Rev. Lett. 92, 092301 (2004).
[77] C. Adler et al., Phys. Rev. C 66, 061901 (2002).
[78] C. Markert, nucl-ex/0404003 (2004).
[79] E. S. Fraga, R. D. Pisarski, and J. Schaffner-Bielich, Phys. Rev. D 63, 121702(R) (2001).