SQM 2006: Theory Summary and Perspectives

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1. Introduction

The investigation of strangeness production in relativistic heavy ion collisions has been proven to be a powerful tool for the study of highly excited nuclear matter, both in terms of the reaction dynamics and in terms of its hadrochemistry [1, 2, 3, 4, 5, 6, 7, 8, 9, 10]. Furthermore, strangeness has been suggested as a signature for the creation of a Quark-Gluon-Plasma (QGP) [11, 12]: in the final state of a heavy-ion collision, subsequent to the formation and decay of a QGP, strangeness has been predicted to be enhanced relative to the strangeness yield in elementary hadron+hadron collisions. It was actually this early work on strangeness as a QGP signature which was instrumental in creating the SQM conference series. However, as witnessed by the work presented at this conference and its predecessors, strangeness production in heavy-ion collisions has developed into an impressively versatile tool for the characterization of QCD matter.

In this write-up I will discuss a selection of key contributions at SQM 2006 which I consider to have a large impact to the current status of the field of strangeness physics or which may have the potential to significantly advance strangeness – or in general flavor physics – in the near future.

I would like to point out that this write-up is not a comprehensive summary covering all theory topics discussed at SQM 2006 – for example, there have been excellent mini-reviews on strangeness in the nuclear astrophysics context by J. Schaffner-Bielich [11] and S. Reddy [12] as well as on the extraction of the nuclear equation of state via near-threshold kaon production by J. Aichelin [13] and the extension of the periodic table into the strangeness and anti-matter sector by W. Greiner [14] – the coverage of any of these would be beyond the scope of this write-up.
2. State of the Art: Lessons and Challenges

2.1. Hadron Yields and Ratios

The impressive success of statistical models in describing the (strange) hadron abundances and ratios at the CERN/SPS and RHIC \cite{4, 5, 8, 9, 10}, and the extracted $\gamma_s$ values close to 1 have led to the common conclusion that chemical freeze-out in ultra-relativistic heavy-ion reactions occurred very close to – or even at hadrochemical equilibrium and that this state most likely has been created by a hadronizing QGP. Several studies, exploring the systematics of hadron ratios as function of collision system and energy have been presented at this meeting, all corroborating the model’s success and previous findings \cite{15}.

However, there remain a number of open questions associated with the success of the statistical model, many of which transcend the scope of the model and need to be addressed either through direct analysis of data beyond yields and spectra or via dynamical non-equilibrium theory approaches based on ab-initio calculations:

- **How did the system achieve chemical equilibrium?** Conventional calculations based on boost-invariant hydrodynamics with rate-equations for quark production \cite{16, 17, 18}, pQCD rate-equations \cite{19} or the Parton Cascade Model \cite{20} all indicate that chemical equilibration (and strangeness saturation) cannot be achieved during realistic life-times of the deconfined phase. It has been suggested (e.g. in \cite{16, 17}) that the system would be driven toward and come close to chemical equilibrium in the subsequent hadronic phase – a scenario which would help to bridge the gap between the calculations indicating insufficient equilibration time in the plasma phase and the copious SPS and RHIC data apparently close to chemical equilibrium at chemical freeze-out. Recent calculations assuming a hadronizing QGP out of chemical equilibrium with subsequent hadronic rescattering have shown that rescattering via binary collisions in the hadronic phase is insufficient to drive the system toward chemical equilibrium before the expansion of the system leads to chemical freeze-out \cite{21}. Currently the most favored approaches for explaining the rapid thermalization of the medium are either centered around turbulent color fields \cite{22, 23, 24, 25, 26} or multi-particle collisions (which may take place either in the deconfined or dense, confined phase of the reaction) \cite{27, 28, 29, 30}. A smoking gun signature for either mechanism remains yet to be established.

- **Do the temperature and chemical potential extracted from the statistical model fits to final state hadron yields or ratios really reflect the thermodynamic state of the system at one particular time during its evolution (i.e. the conditions at chemical freeze-out) or are they rather the result of a superposition of different states, due to individual hadron species decoupling continuously from the system (as would be expected from their different mean free paths)?** The latter view is supported by a transport model analysis of the time-evolution of the temperature and chemical potential in the central cell of a heavy-ion reaction \cite{31}. However,
one would expect a fairly large chi-squared for the single source statistical model fit in case of a continuous decoupling scenario, which is not observed in the data. It should also be noted that the use of $4\pi$ integrated yields at SPS energies creates ambiguities in the fit, since many hadron ratios have been shown to exhibit a strong rapidity dependence at SPS (and lower) beam energies [32, 33]. True progress on this issue can only be achieved if a model-independent method is found to determine the chemical decoupling time of individual hadron species.

- taking the extracted temperature and chemical potential values for SPS and RHIC at face value, the system chemically decouples very close to $T_C$ – at a temperature at which the properties of hadrons (e.g. their mass and/or width) could still be substantially modified by the temperature and density of the medium (see e.g. the spectral function of the $\rho$ presented at this meeting by A. Foerster [34]). Is the use of hadronic vacuum masses and widths in the statistical model the right approach or should their temperature- and density dependent medium modifications be taken into account in the model fits [35, 36]?

Overall, to this day the physics mechanisms and driving forces behind the impressive success of the statistical model are not well understood – finding an answer to the question why the statistical model performs so well is a challenge the theory community must address with renewed vigor in the near future.

### 2.2. Excitation Functions

Over the past decade the experimental programs at the AGS, SPS and RHIC have resulted in a wealth of data which can be compiled into excitation functions of yields, ratios, flow coefficients, system sizes etc.. However, the hope of finding some kind of sharp discontinuity as a signal of a phase transition has not come to fruition. This may be partially due to many QGP signatures predicted to exhibit a sharp discontinuity or local extremum relying heavily on the deconfinement phase-transition being of first order with a long-lived mixed phase – an assumption which is at odds with recent findings from Lattice QCD, predicting the phase transition to be a continuous crossover in the RHIC and SPS domain. Even in the case of a 1st order phase-transition the expectation of sharp features in excitation functions may have been somewhat naive: sharp discontinuities are predicted as function of temperature, which does not have a one-to-one relation to beam energy. Realistic temperature and density profiles as well as corona effects (see e.g. the contribution by K. Werner [37]) would lead to a smearing out of the features associated with deconfinement and thus would make for a smooth variation in the excitation function.

Spurred by recent lattice-gauge theory calculations of QCD at finite baryon density [38] there has been renewed hope that an excitation function in incident-beam energy might yield interesting results if one can create a QCD medium close to the tri-critical point. It has been found that the net-quark susceptibility near the tri-critical point diverges, which would be experimentally accessible via charge fluctuation measurements.
2.3. Transport Theory

Relativistic Fluid Dynamics (RFD, see e.g. [41, 42, 43]) is ideally suited for the high-density phase of heavy-ion reactions at RHIC, but breaks down in the later, dilute, stages of the reaction when the mean free paths of the hadrons become large and flavor degrees of freedom are important. The biggest advantage of RFD is that it directly incorporates an equation of state as input and thus is so far the only dynamical model in which a phase transition can explicitly be incorporated. In the ideal fluid approximation (i.e. neglecting off-equilibrium effects) and once the initial conditions for the calculation have been fixed, the EoS is the only input to the equations of motion and relates directly to properties of the matter under consideration. Ideally, either the initial conditions or the EoS should be determined beforehand by an ab-initio calculation (e.g. for the EoS via a lattice-gauge calculation), in which case a fit to the data would allow for the determination of the remaining quantity.

RFD has been extremely successful in describing single particle spectra and collective flow effects at RHIC [44, 45, 46], even though no hydrodynamical model implementation has so far attempted to address the entire array of available data in a single consistent calculation (a forthcoming publication aims to remedy this situation [48]). The shape of the spectra as well as the transverse momentum dependence of the elliptic flow for minimum bias data are generally reproduced nicely. However, more specific centrality bins pose a problem for the elliptic flow calculations. One should also bear in mind that hydrodynamical calculations which assume a standard chemical equilibrium hadron gas equation of state below $T_c$ (which implies simultaneous chemical and thermal freeze-out) are unable to fit the measured particle yields simultaneously with the spectra, since statistical models show chemical freeze-out to occur around $T=170$ MeV, whereas the shape of the spectra requires a hydrodynamic evolution to $T=110$ MeV. One method to deal with the separation of chemical and thermal freeze-out is the partial chemical equilibrium model (PCE) [49, 50, 51]. Below a chemical freeze-out temperature $T_{ch}$ one introduces a chemical potential for each hadron whose yield is supposed to be frozen out at $T_{ch}$. The PCE approach can account for the proper normalization of the spectra, however, it fails to reproduce the transverse momentum and mass dependence of the elliptic flow. More importantly for the strangeness sector, ideal RFD calculations lack the capability of dealing with the flavor dependence of hadronic cross sections and a possible flavor-dependent sequential freeze-out: recent experimental results suggest that at thermal freeze-out multistrange baryons exhibit less transverse flow and a higher temperature closer to the chemical freeze-out temperature compared to non- or single-strange baryons [52, 53]. This behavior can be understood in terms of the flavor dependence of the hadronic cross section, which decreases with increasing strangeness content of the hadron. The reduced cross section of multi-strange baryons leads to a decoupling from the hadronic medium at an earlier stage of the
reaction, allowing them to provide information on the properties of the hadronizing QGP less distorted by hadronic final state interactions [54, 55].

The reach of RFD can be extended and the problem having to terminate the calculation at a single flavor-independent and fixed freeze-out temperature can be overcome by combining the RFD calculation with a microscopic hadronic cascade model – this kind of hybrid approach (dubbed hydro plus micro) was pioneered in [58] and has subsequently been adopted by other groups [59, 61, 47, 48]. Its key advantages are that the freeze-out occurs naturally as a result of the microscopic evolution and that flavor degrees of freedom are treated explicitly through the hadronic cross sections of the microscopic transport. Due to the Boltzmann equation being the basis of the microscopic calculation in the hadronic phase, viscous corrections for the hadronic phase are by default included in the approach – the full treatment of viscosity in the deconfined phase in a 3D hydrodynamic calculation remains a challenge for the future.

Such hybrid macro/micro transport calculations are to date the most successful approaches for describing the soft physics at RHIC [62]. First implementations of these hybrid approaches were restricted to 1+1 [58] and 2+1 [59] dimensions in the hydrodynamic component of the model. However, recently new state of the art fully three-dimensional ideal hydrodynamics models have become available and are now being incorporated into the hydro+micro framework [47, 60, 61, 48].

The latest model implementation has been shown at this conference [48] to be able to excellently reproduce the measured $P_T$ spectra and $\langle P_T \rangle$ values of hyperons and multi-strange baryons at RHIC (see left frame of figure 1). A collision number analysis for these hadron species based on this calculation [48] confirms the early and sequential freeze-out findings described in [54, 55].

Marcus Bleicher discussed elliptic flow at RHIC in the framework of the microscopic hadron/string UrQMD and RQMD transport models [65, 66]. Both models reach only 60% of the absolute magnitude of the measured $v_2$, due to the string degrees of freedom being used to describe the initial phase in these models generating insufficient pressure. Interestingly, the model calculations exhibit features reminiscent of constituent quark number scaling in the intermediate $p_T$ range between 2 and 6 GeV/c. This is due to the use of the Additive Quark Model for most hadron-hadron cross section implemented in the model, thus having the interaction scale with the number of constituent quarks. At closer inspection, however, one finds a flavor ordering in the elliptic flow generated by the microscopic transport models: $v_2(N) > v_2(Y) > v_2(\Xi) > v_2(\Omega)$, due to the strange quark being assigned a smaller interaction cross section than the $u$ and $d$ quarks. This structure is not observed in the data, which rather indicate the build-up of elliptic flow independent of the quark flavor prior to hadronization.

2.4. Hadronization via Parton Recombination

The detailed understanding of hadronization plays a crucial role for isolating signatures sensitive to the QGP evolution and properties of the system from those which are
Figure 1. Left: $P_T$ spectra for $\Lambda$, $\Xi$, $\Omega$ in central collisions calculated in the hybrid 3D-hydro+UrQMD approach compared to STAR data. Right: elliptic flow of (multi-)strange hadrons measured by the STAR collaboration [56, 57]. Note the $\phi$ meson following the kaon elliptic flow.

dominated by the later reaction stages. One of the theoretical milestones of the first several years of the RHIC program was the development of the recombination plus fragmentation model as the standard model of hadronization for matter at intermediate and high transverse momenta: at the center of the recombination + fragmentation model is the realization that hadron production at momenta of a few GeV/c in an environment with a high density of partons occurs by recombination, rather than fragmentation, of partons. It is found that recombination always dominates over fragmentation for an exponentially falling parton spectrum, but that fragmentation wins out eventually, when the spectrum takes the form of a power law.

Among the RHIC discoveries that prompted the development of the recombination models of hadronization, is that the amount of suppression of hadron production at intermediate $p_T$ compared to the scaled proton-proton baseline seems to depend on the hadron species. In fact, in the production of protons and antiprotons between 2 and 4 GeV/c the suppression seems to be completely absent. Generally, pions and kaons appear to suffer from a strong energy loss while baryons and antibaryons do not. Two stunning experimental facts exemplify this [63, 64]. First, the ratio of protons over positively charged pions is equal or above one for $p_T > 1.5$ GeV/c and is approximately constant up to 4 GeV/c. Second, the nuclear suppression factor $R_{AA}$ below 4 GeV/c is close to one for protons and lambdas [67], while it is about 0.3 for pions.

The recombination approach [68, 69, 70, 71] has been able to account for the above baryon/meson differences. Additionally, it was observed that the elliptic flow pattern of different hadron species can be explained by a simple recombination mechanism [72, 73, 74, 75]. The anisotropies $v_2$, for the different hadrons are compatible with
a universal function \( v_2(p_T) \) in the parton phase, related to the hadronic flow by factors of two and three depending on the number of valence quarks \[76, 77\].

One of the perceived weaknesses of the recombination+fragmentation approach was that its development was triggered by the experimental observation of the meson/baryon anomalies in the RHIC data and therefore its success in explaining these features occurred after the fact. However, the recombination approach was able to predict the the elliptic flow \( v_2 \) of multi-strange hadrons such as the \( \phi \), \( \Xi \) and \( \Omega \) as a function of transverse momentum \[78\]: the measurement of \( v_2 \) for the \( \phi \) and \( \Omega \) allows for the unambiguous distinction between parton recombination and statistical hadro-chemistry to be the dominant process in hadronization at intermediate transverse momenta, since e.g. the \( \phi \) meson has approximately the mass of a nucleon, but the valence-quark content of a meson. In a hydrodynamic picture with the hadron mass as the guiding scale the \( \phi \) would follow the systematics of the nucleon whereas in the recombination picture it would follow the behavior of the pions and kaons. Data on the elliptic flow of the \( \phi \) meson and \( \Omega \) baryon presented at this conference, e.g. in the talk by S.L. Blyth – see the right frame of figure 1 – clearly exhibit constituent quark scaling and therefore impressively confirm the physics of the recombination model and demonstrate its predictive power.

3. New Directions

3.1. Quarks in the Color Glass

Heavy-ion collisions at ultra-relativistic energies can be described by the collision of two coherent sheets of high energy density gluonic fields (commonly referred to as Color Glass Condensate) \[79, 80, 81\]. Since the physical density of gluons becomes large, their typical separation is small, implying a small value for \( \alpha_s \). Furthermore, these highly coherent gluons saturate the phase space up to the maximal occupation number \( \sim \frac{1}{\alpha_s} \). Due to the weak coupling, it is possible to describe this system from first principles in QCD.

The Color Glass Condensate (CGC) has been suggested to describe the initial state of gold nuclei in RHIC collisions. While there is a broad consensus that at sufficiently high beam energies the saturation physics of the CGC should dominate the initial state, it is still a matter of debate whether these conditions are actually fulfilled at RHIC. Nonetheless, the further development of the CGC picture for heavy-ion collisions is of great theoretical importance. Original work on the CGC focused on the description of the initial state as gluonic field. Due to the large gluon densities at low Bjorken-x, these are thought to dominate the dynamics of the initial state. However, the lack of treatment and consideration of quark and anti-quark production makes it difficult to connect the CGC to experimental data, e.g. in the strangeness sector.

T. Lappi presented a numerical integration of the Dirac equation in order to calculate the number of quark-antiquark pairs initially produced in the classical color
fields of colliding ultra-relativistic nuclei (the resulting state of high energy and density matter being termed the \textit{glasma}) \cite{2}. While the number of $q\bar{q}$ pairs is parametrically suppressed in the coupling constant, he found that in the CGC their production rate is comparable to the thermal ratio of gluons/pairs $= 9N_f/32$. After isotropization one thus would end up with a quark-gluon plasma in chemical equilibrium. This finding is of great significance since it provides a link between the CGC initial state and an abundance of flavor-centric data indicating the formation of a QGP in thermal and chemical equilibrium.

### 3.2. Conserved Charge Correlations

The elementary electric charge, strangeness and baryon number carried by the QGP degrees of freedom differs significantly from those in a gas of hadrons. This observation does not only apply to the elementary charges themselves, but also to the size of the average fluctuations of net baryon number, strangeness and electric charge in a finite volume. While hadronization and confinement prohibit us from directly observing the fractional electric charge and baryon number and strangeness of the QGP degrees of freedom, these event by event fluctuations may under certain conditions survive hadronization and subsequent hadronic rescattering and thus serve as indicators of the existence of a QGP \cite{3, 4}.

Surprisingly, all experimental analysis of net charge fluctuations for SPS and RHIC data agree with the hadron gas prediction, giving no indication at all about a possible deconfined phase. Since many other measurements are compatible with the assumption of deconfinement, it is tempting to speculate that the dynamics of hadronization and/or hadronic final state interactions strongly affect the charge fluctuation observable and mask the fluctuations generated in the deconfined phase.

A more robust approach of exploiting the difference in conserved elementary charges between QGP and HG degrees of freedom lies in the area of correlations: in particular, the correlation between baryon number $B$ and strangeness $S$ may elucidate the microscopic structure of the QGP and the nature of its degrees of freedom \cite{5, 6, 7}. Strangeness is carried exclusively by $s$ and $\bar{s}$ quarks which carry baryon number as well: $B_s = -\frac{1}{3}S$, whereas in a hadron gas the relation between $B$ and $S$ is much less intimate. This $C_{BS}$ correlation can be directly extracted from Lattice QCD calculations and is experimentally accessible via event-by-event fluctuations:

$$C_{BS} \equiv -3\frac{\sigma_{BS}}{\sigma_S^2} = -3\frac{\langle BS \rangle - \langle B \rangle \langle S \rangle}{\langle S^2 \rangle - \langle S \rangle^2} = -3\frac{\langle BS \rangle}{\langle S^2 \rangle}$$

(1)

For a system of hadrons $C_{BS}$ can be formulated in terms of multiplicity variances $\sigma_k^2 \equiv \langle n_k^2 \rangle - \langle n_k \rangle^2 \approx \langle n_k \rangle$ for the respective hadron species $k$:

$$C_{BS} = -3\frac{\sum_k \sigma_k^2 B_k S_k}{\sum_k \sigma_k^2 S_k^2} \approx -3\frac{\sum_k \langle n_k \rangle B_k S_k}{\sum_k \langle n_k \rangle S_k^2}$$

(2)
For an ideal gas of hadrons at $T = 170$ MeV and zero chemical potential $\mu_B = 0$ one finds $C_{BS} = 0.66$.

The $C_{BS}$ correlator can also be expressed in terms of susceptibilities, $C_{BS} = -3\chi_{BS}/\chi_{SS}$ which are second derivatives of the free energy with respect to the chemical potential and can be directly be calculated from lattice QCD:

$$\chi_{BS} = -\frac{1}{V} \frac{\partial^2 F}{\partial \mu_B \partial \mu_S}, \quad \chi_{SS} = -\frac{1}{V} \frac{\partial^2 F}{\partial \mu_S^2}. \quad (3)$$

Using values of lattice susceptibilities extracted at $T = 1.5T_C$ from lattice QCD on finds $C_{BS} \approx 1$ indicating that the degrees of freedom of a QGP carry baryon number and strangeness of individual quarks and suggesting that the quark flavors are uncorrelated as in an ideal QGP (see left frame of figure 2). Note that the presence of pure gluon clusters cannot be ruled out by this diagnostic. This finding is of great importance, since the evaluation of $C_{BS}$ for different models of the QGP structure allows for the direct verification or falsification of these models when compared to data and lattice QCD.

### 3.3. Anomalous Viscosity

As stated earlier, measurements of the anisotropic collective flow of hadrons emitted in non-central collisions of heavy nuclei at the Relativistic Heavy Ion Collider (RHIC) are
in remarkably good agreement with the predictions of ideal relativistic fluid dynamics. In order to describe the data, calculations need to assume that the matter formed in the nuclear collision reaches thermal equilibrium within a time \( \tau_i \approx 1 \text{ fm/c} \) and then expands with a very small shear viscosity \( \eta \ll s \), where \( s \) is the entropy density. The comparison between data and calculations indicates that the viscosity of the matter cannot be much larger than the postulated lower bound \( \eta_{\text{min}} = s/4\pi \), which is reached in certain strongly coupled supersymmetric gauge theories.

This result is nontrivial because the shear viscosity of a weakly coupled, perturbative quark-gluon plasma is not small. In fact, the perturbative result for the shear viscosity, in leading logarithmic approximation, is

\[
\eta_C = \frac{d_f T^3}{g^4 \ln g^{-1}},
\]

where \( d_f \approx O(100) \) is a numerically determined constant that weakly depends on the number of quark flavors \( n_f \). The result, as well as the finding that numerical solutions of the Boltzmann equation exhibit fluid dynamical behavior only when the cross section between gluons is artificially increased by a factor ten or more, have invited speculations that the matter produced at RHIC is a strongly coupled quark-gluon plasma (sQGP). The possible microscopic structure of such a state is not well understood at present.

However, as Berndt Müller pointed out in his talk, there exists an alternative mechanism that may be responsible for a small viscosity of a weakly coupled, but expanding quark-gluon plasma. This mechanism is based on the theory of particle transport in turbulent plasmas. Such plasmas are characterized by strongly excited random field modes in certain regimes of instability, which coherently scatter the charged particles and thus reduce the rate of momentum transport. The scattering by turbulent fields in electromagnetic plasmas is known to greatly increase the energy loss of charged particles and reduce the heat conductivity and the viscosity of the plasma. Following Abe and Niu, the contribution from turbulent fields to transport coefficients was called “anomalous”.

The sufficient condition for the spontaneous formation of turbulent, partially coherent fields is the presence of instabilities in the gauge field due to the presence of charged particles. This condition is met in electromagnetic plasmas with an anisotropic momentum distribution of the charged particles, and it is known to be satisfied in quark-gluon plasmas with an anisotropic momentum distribution of thermal partons.

The additional contribution to the viscosity, \( \eta_B \), induced by the turbulent fields decreases with increasing strength of the fields. Since the amplitude of the turbulent fields grows with the magnitude of the momentum anisotropy, a large anisotropy will lead to a small value of \( \eta_B \). Because the relaxation rates due to different processes are additive, the total viscosity is given by

\[
\eta^{-1} = \eta_B^{-1} + \eta_C^{-1}.
\]
This equation implies that $\eta_B$ dominates the total shear viscosity, if it is smaller than $\eta_C$. In that limit, the anomalous mechanism exhibits a stable equilibrium in which the viscosity regulates itself: The anisotropy grows with $\eta$, but an increased anisotropy tends to suppress $\eta_B$. The result is that in the weak coupling limit, the anomalous viscosity is much smaller than the viscosity due to collisions among thermal partons. By reducing the shear viscosity of a weakly coupled, but expanding quark-gluon plasma, this mechanism could possibly explain the observations of the RHIC experiments without the assumption of a strongly coupled plasma state.

4. From SQM to F(lavor)QM

4.1. Heavy-Quark Production, Diffusion and Energy Loss

The calculation of heavy quark production in the framework of pQCD is fairly well established – recent improvements having been reported by I. Vitev at this meeting [107] (see also [106] and references therein). Among the novelties presented was the work by K. Tuchin, who discussed heavy quark production in High Parton Density QCD in a quasi-classical approximation, including low-$x$ quantum evolution, as well as heavy-quark production based on the effect of pair production in external fields [108].

For moderate transverse momenta ($p_T \lesssim$ a few GeV/$c$) the energy loss of heavy quarks is thought to be dominated by inelastic collisions with medium constituents, rather than gluon radiation [109, 110], because the heavy quarks are not ultra-relativistic and gluon radiation is suppressed by the so-called dead-cone effect [111]. These findings were highlighted in a comprehensive analysis of the non-photonic electron nuclear suppression factors measured at RHIC presented by M. Djordjevic [112, 113]. The analysis clearly indicates the importance of including charm and bottom quarks in the calculation and taking their radiative as well as collisional energy loss into account.

The importance of the collisional energy loss contribution and the possibility of heavy quarks actually thermalizing in the medium provides an opportunity to utilize them as probes for the transport coefficients of the QCD medium. Since collisional energy loss occurs in many small steps, the motion of a heavy quark can thus be described by a Fokker-Planck or, equivalently, Langevin equation. Several studies based on such an approach have recently been done [114, 115, 116]. The most detailed and extensive one of these was performed by Moore and Teaney [114], who derived an expression for the diffusion coefficient $D$ in the framework of hard-thermal loop (HTL) improved perturbation theory and discussed the limitations of the Fokker-Planck approach. These authors also studied the resulting phenomenology of heavy quark transport in dense matter created by a heavy ion collision, using a boost invariant hydrodynamical model with an ideal equation of state. At SQM 2006, heavy quark diffusion calculations were discussed by P.B. Gossiaux [116, 117] as well as by R. Rapp [115, 118]. Whereas the philosophy of [116, 117] is to determine the transport coefficients from a comparison to data, the approach of [115, 118] is based on introducing a resonant charm – light quark
interaction and then calculating the drag- and diffusion coefficients for the Langevin evolution on the basis of that interaction. A comparison to data would thus yield information on the microscopic interaction between charm and the light quark species (see right frame of figure 2).

4.2. Charmonium Spectral Functions

In 1986 Matsui and Satz proposed [119] that the suppression of heavy quarkonia-mesons could provide one of the signatures for deconfinement in QCD at high temperatures. The idea was based on an analogy with the well known Mott transition in condensed matter systems. At high densities, Debye screening in a quark-gluon plasma reduces the range of the attractive force between heavy quarks and antiquarks, and above some critical density screening prevents the formation of bound states. The larger bound states were expected to dissolve before the smaller ones as the temperature of the system increases. The $\psi'$ and $\chi_c$ states were thus expected to become unbound just above $T_c$, while the smaller $\psi$ state would only dissolve above $\approx 1.2T_c$.

However, QCD lattice gauge-theory calculations of charmonium correlators in recent years have necessitated a revision in our understanding of the dissociation of charmonium states and what it implies for the properties of the surrounding medium: as discussed by P. Petreczky at this meeting, the 1S charmonia states ($J/\psi$, $\psi'$ and $\eta_C$) survive to unexpectedly high temperatures above 1.5 $T_C$ and only the 1P states (i.e. the $\chi_C$) dissolves around 1.2 $T_C$ [120, 121]. It is therefore questionable whether the observed charmonium suppression in Pb+Pb collisions at the SPS [122] is truly a smoking gun signature for deconfinement or rather the result hadronic dissociation (see e.g. [123]). Moving on to RHIC energies, the situation will become even more complicated: achievable temperatures may be above the threshold for charmonium dissociation (which is unlikely in the SPS case), however, the additional suppression may be compensated through novel production mechanisms. The multiplicity of produced charm and anti-charm quarks per event at RHIC is sufficiently large that charmonium-regeneration via parton recombination may occur [124, 125], giving rise to a fairly flat behavior of the observed charmonium yield as a function of beam energy.

Another exciting prospect is the measurement of charmonium elliptic flow, first shown in form of preliminary data by NA60 here at this conference [34]. The measurement of significant charmonium elliptic flow would indicate charm thermalization and charmonium regeneration via parton recombination. In addition it has been shown [118], that this observable exhibits a strong sensitivity to the in-medium interaction of charm quarks.

5. Outlook

The field of strangeness in heavy-ion collisions – and with it this conference series – has developed tremendously over the past 20 years. While strangeness as a QGP signature
was instrumental in creating the SQM conference series, strangeness production in heavy-ion collisions has become an impressively versatile tool for the characterization of all aspects of confined and deconfined QCD matter, as witnessed by the work presented at this conference. The future of the field will most likely bring the generalization of the concepts and lessons learned from strangeness into the heavy-quark sector at higher incident beam energies. The application of flavor-dominated physics phenomena as probes of the the hot and dense QCD medium will be an exciting field for many years to come.

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References

[1] J. Rafelski and B. Müller, Phys. Rev. Lett. 48, 1066(1982)
[2] P. Koch, B. Müller, J. Rafelski, Phys. Rep. 142, 167 (1986).
[3] P. Koch, B. Müller, H. Stöcker, W. Greiner, Mod. Phys. Lett. A3, 737 (1988).
[4] P. Braun-Munzinger, J. Stachel, J. P. Wessels and N. Xu, Phys. Lett. B 344, 43 (1995).
[5] J. Letessier, A. Tounsi and J. Rafelski, Phys. Lett. B 389, 586 (1996).
[6] S. E. Vance and M. Gyulassy, Phys. Rev. Lett. 83, 1735 (1999).
[7] S. Soff, S. A. Bass, M. Bleicher, L. Bravina, E. Zabrodin, H. Stocker and W. Greiner, Phys. Lett. B 471, 89 (1999).
[8] P. Braun-Munzinger, I. Heppe and J. Stachel, Phys. Lett. B 465, 15 (1999).
[9] J. Rafelski, J. Letessier and G. Torrieri, Phys. Rev. C 64, 054907 (2001) [Erratum-ibid. C 65, 069902 (2002)].
[10] P. Braun-Munzinger, D. Magestro, K. Redlich and J. Stachel, Phys. Lett. B 518, 41 (2001).
[11] J. Schaffner-Bielich, these proceedings.
[12] S. Reddy, these proceedings.
[13] J. Aichelin, these proceedings.
[14] W. Greiner, these proceedings.
[15] P. Braun-Munzinger, these proceedings;
    J. Cleymans, these proceedings;
    talk by I. Kraus at SQM 2006.
[16] J. Kapusta and A. Mekjian. Phys. Rev. D33 (1986) 1304.
[17] T. Matsui, B. Svetisky and L. D. McLerran. Phys. Rev. D34 (1986) 2047.
[18] D.M. Elliott and D.H. Rischke. Nucl. Phys. A671 (2000) 583.
[19] T. Biro, E. van Doorn, B. Müller, M. H. Thoma and X. N. Wang. Phys. Rev. C48 (1993) 1275.
[20] K. Geiger and J. I. Kapusta. Phys. Rev. D47 (1993) 4905.
[21] S. A. Bass, P. Danielewicz, S. Pratt and A. Dumitru, J. Phys. G 27, 635 (2001).
[22] S. Mrowczynski, Phys. Lett. B 214 (1988) 587.
[23] S. Mrowczynski, Phys. Lett. B 314 (1993) 118.
[24] P. Romatschke and M. Strickland. Phys. Rev. D 68, 036004 (2003)
[25] J. Randrup and S. Mrowczynski, Phys. Rev. C 68, 034909 (2003)
[26] P. Arnold, J. Lenaghan and G. D. Moore, JHEP 0308, 002 (2003)
[27] Z. Xu and C. Greiner, Phys. Rev. C 71, 064901 (2005) arXiv:hep-ph/0406278.
[28] R. Rapp and E. V. Shuryak, collisions,” Phys. Rev. Lett. 86, 2980 (2001).
[29] C. Greiner and S. Leupold, J. Phys. G 27, L95 (2001).
[30] P. Braun-Munzinger, J. Stachel and C. Wetterich, Phys. Lett. B 596, 61 (2004) arXiv:nucl-th/0311005.
[31] L. V. Bravina et al., Phys. Rev. C 63, 064902 (2001) arXiv:hep-ph/0010172.
[32] S. A. Bass et al., Phys. Rev. Lett. 81, 4092 (1998) arXiv:nucl-th/9711032.
[33] J. Sollfrank, U. W. Heinz, H. Sorge and N. Xu, Phys. Rev. C 59, 1637 (1999) arXiv:nucl-th/9811011.
[34] A. Foerster, these proceedings.
[35] M. Michalec, W. Florkowski and W. Broniowski, Phys. Lett. B 520, 213 (2001) arXiv:nucl-th/0103029.
[36] D. Zschiesche, S. Schramm, J. Schaffner-Bielich, H. Stoecker and W. Greiner, Phys. Lett. B 547, 7 (2002) arXiv:nucl-th/0209022.
[37] K. Werner, these proceedings.
[38] R. Gavai, these proceedings.
[39] S. Ejiri, F. Karsch and K. Redlich, Phys. Rev. C 63, 064902 (2001) arXiv:hep-ph/0010172.
[40] K. Redlich, these proceedings.
[41] J. D. Bjorken, Phys. Rev. D 27, 140 (1983).
[42] R. B. Clare and D. Strottman, Phys. Rept. 141, 177 (1986).
[43] A. Dumitru and D. H. Rischke, Phys. Rev. C 59, 354 (1999) arXiv:nucl-th/9806003.
[44] P. F. Kolb and U. W. Heinz, arXiv:nucl-th/0305084.
[45] P. Huovinen, arXiv:nucl-th/0305064.
[46] T. Hirano and K. Tsuda, Nucl. Phys. A 715, 821 (2003) arXiv:nucl-th/0208068.
[47] C. Nonaka and S. A. Bass, arXiv:nucl-th/0510038.
[48] C. Nonaka and S. A. Bass, arXiv:nucl-th/0607018.
[49] N. Arbex, F. Grassi, Y. Hama and O. Socolowski, Phys. Rev. C 64, 064906 (2001).
[50] T. Hirano and K. Tsuda, Phys. Rev. C 66, 054905 (2002) arXiv:nucl-th/0205043.
[51] P. F. Kolb and R. Rapp, Phys. Rev. C 67, 044903 (2003) arXiv:hep-ph/0210222.
[52] M. Estienne (for the STAR Collaboration), J. Phys. G31, S873 (2005).
[53] J. Adams et al. (STAR Collaboration), Phys. Rev. Lett. 92, 182301 (2004).
[54] H. van Hecke, H. Sorge and N. Xu, Phys. Rev. Lett. 81, 5764 (1998) arXiv:nucl-th/9804035.
[55] A. Dumitru, S. A. Bass, M. Bleicher, H. Stoecker and W. Greiner, Phys. Lett. B 460, 411 (1999) arXiv:nucl-th/9901046.
[56] S.L. Blyth, these proceedings.
[57] J. Adams et al. [STAR Collaboration], Phys. Rev. Lett. 95, 122301 (2005) arXiv:nucl-ex/0504022.
[58] S. A. Bass and A. Dumitru, Phys. Rev. C 61, 064909 (2000) arXiv:nucl-th/0001033.
[59] D. Teaney, J. Lauret and E. V. Shuryak, Phys. Rev. Lett. 86, 4783 (2001) arXiv:nucl-th/0111058.
[60] T. Hirano, U. W. Heinz, D. Kharzeev, R. Lacey and Y. Nara, Phys. Lett. B 636, 299 (2006) arXiv:nucl-th/0511046.
[61] T. Hirano, arXiv:nucl-th/0510005.
[62] T. Hirano and M. Gyulassy, arXiv:nucl-th/0506049.
[63] K. Adcox et al. [PHENIX Collaboration], Phys. Rev. Lett. 88, 242301 (2002) arXiv:nucl-ex/0112006.
[64] S. S. Adler et al. [PHENIX Collaboration], Phys. Rev. Lett. 91, 172301 (2003) arXiv:nucl-ex/0305036.
[65] Y. Lu et al., arXiv:nucl-th/0602009.
[66] M. Bleicher, these proceedings.
[67] J. Adams et al. [STAR Collaboration], Phys. Rev. Lett. 92, 052302 (2004) arXiv:nucl-ex/0306007.
[68] R. J. Fries, B. Müller, C. Nonaka and S. A. Bass, Phys. Rev. Lett. 90, 202303 (2003).
[69] R. J. Fries, B. Muller, C. Nonaka and S. A. Bass, Phys. Rev. C 68, 044902 (2003).
[70] V. Greco, C. M. Ko and P. Levai, Phys. Rev. Lett. 90, 202302 (2003).
[71] R. C. Hwa and C. B. Yang, [arXiv:nucl-th/0211010].
[72] Z. W. Lin and C. M. Ko, Phys. Rev. Lett. 89, 202302 (2002).
[73] S. A. Voloshin, Nucl. Phys. A 715, 379c (2003).
[74] D. Molnár and S. A. Voloshin, Phys. Rev. Lett. 91, 092301 (2003) [arXiv:nucl-th/0302014].
[75] Z. W. Lin and C. M. Ko, Phys. Rev. Lett. 89, 202302 (2002).
[76] S. A. Voloshin, Nucl. Phys. A 715, 379c (2003).
[77] D. Molnár and S. A. Voloshin, Phys. Rev. Lett. 91, 092301 (2003) [arXiv:nucl-th/0302014].
[78] Z. Lin and D. Molnár, Phys. Rev. C 68, 044901 (2003) [arXiv:nucl-th/0304045].
[79] P. Sorensen for the STAR Collaboration, J. Phys. G 30, S217 (2004) [arXiv:nucl-ex/0305008].
[80] S. S. Adler et al. [PHENIX Collaboration], Phys. Rev. Lett. 91, 182301 (2003) [arXiv:nucl-ex/0305013].
[81] C. Nonaka, R. J. Fries and S. A. Bass, Phys. Lett. B 583, 73 (2004) [arXiv:nucl-th/0308051].
[82] F. Gelis, K. Kajantie and T. Lappi, Phys. Rev. Lett. 96, 032304 (2006) [arXiv:hep-ph/0508229].
[83] M. Asakawa, U. W. Heinz and B. Muller, Phys. Rev. Lett. 85, 2072 (2000) [arXiv:hep-ph/0003169].
[84] S. Jeon and V. Koch, Phys. Rev. Lett. 85, 2076 (2000) [arXiv:hep-ph/0003168].
[85] V. Koch, A. Majumder and J. Randrup, Phys. Rev. Lett. 95, 182301 (2005) [arXiv:nucl-ex/0505052].
[86] A. Majumder and B. Muller, [arXiv:nucl-th/0605079].
[87] A. Majumder, these proceedings.
[88] R. V. Gavai and S. Gupta, Phys. Rev. D 73, 014004 (2006) [arXiv:hep-lat/0510044].
[89] U. W. Heinz and P. F. Kolb, Nucl. Phys. A 702, 269 (2002).
[90] D. Teaney, Phys. Rev. C 68, 034913 (2003).
[91] P. K. Kovtun, D. T. Son and A. O. Starinets, Phys. Rev. Lett. 94, 111601 (2005).
[92] G. Policastro, D. T. Son and A. O. Starinets, Phys. Rev. Lett. 87, 081601 (2001).
[93] P. Arnold, G. D. Moore and L. G. Yaffe, JHEP 0011, 001 (2000); ibid. 0305, 051 (2003).
[94] D. Molnar and M. Gyulassy, Nucl. Phys. A 697, 495 (2002) [Erratum ibid. 703, 893 (2002)].
[95] E. V. Shuryak and I. Zahed, Phys. Rev. D 70, 054507 (2004).
[96] J. Liao and E. V. Shuryak, [arXiv:hep-ph/0510110].
[97] T. H. Dupree, Phys. Fluids 9, 1773 (1966).
[98] M. Asakawa, S. A. Bass and B. Muller, [arXiv:hep-ph/0603092].
[99] T. H. Dupree, Phys. Fluids 11, 2680 (1968).
[100] T. Okada and K. Niu, J. Plasma Phys. 23, 423 (1980).
[101] R. C. Malone, R. L. McCrory, and R. L. Morse, Phys. Rev. Lett. 34, 721 (1975).
[102] T. Okada, T. Yabe, and K. Niu, J. Plasma Phys. 20, 405 (1978).
[103] T. Abe and K. Niu, J. Phys. Soc. Japan 49, 717 (1980).
[104] T. Abe and K. Niu, J. Phys. Soc. Japan 49, 725 (1980).
[105] E. S. Weibel, Phys. Rev. Lett. 2, 83 (1959).
[106] R. Vogt, Eur. Phys. J. C 43, 113 (2005) [arXiv:hep-ph/0412302].
[107] I. Vitev, these proceedings.
[108] K. Tuchin, these proceedings.
[109] M. G. Mustafa and M. H. Thoma, Acta Phys. Hung. A 22, 93 (2005) [arXiv:hep-ph/0311168].
[110] M. G. Mustafa, Phys. Rev. C 72, 014905 (2005) [arXiv:hep-ph/0412402].
[111] Y. L. Dokshitzer, V. A. Khoze and S. I. Troian, J. Phys. G 17, 1602 (1991).
[112] M. Djordjevic, these proceedings.
[113] S. Wicks, W. Horowitz, M. Djordjevic and M. Gyulassy, [arXiv:nucl-th/0512076].
[114] G. D. Moore and D. Teaney, Phys. Rev. C 71, 064904 (2005) [arXiv:hep-ph/0412346].
[115] H. van Hees and R. Rapp, Phys. Rev. C 71, 034907 (2005) [arXiv:nucl-th/0412015].
[116] P. B. Gossiaux, V. Guiho and J. Aichelin, J. Phys. G 31, S1079 (2005) [arXiv:hep-ph/0411324].
[117] P.B. Gossiaux, these proceedings.
[118] R. Rapp and H. van Hees, arXiv:hep-ph/0606117.
[119] T. Matsui and H. Satz, Phys. Lett. B 178, 416 (1986).
[120] P. Petreczky, these proceedings.
[121] A. Jakovac, P. Petreczky, K. Petrov and A. Velytsky, arXiv:hep-lat/0603005.
[122] M. C. Abreu et al. [NA50 Collaboration], Phys. Lett. B 477, 28 (2000).
[123] C. Spieles, R. Vogt, L. Gerland, S. A. Bass, M. Bleicher, H. Stoecker and W. Greiner, Phys. Rev. C 60, 054901 (1999) arXiv:hep-ph/9902337.
[124] R. L. Thews, M. Schroedter and J. Rafelski, Phys. Rev. C 63, 054905 (2001) arXiv:hep-ph/0007323.
[125] L. Grandchamp and R. Rapp, Phys. Lett. B 523, 60 (2001) arXiv:hep-ph/0103124.