HEAVY QUARK CHEMICAL POTENTIAL AS PROBE OF THE PHASE DIAGRAM OF NUCLEAR MATTER

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We study the temperature dependence of the strange and charm quark chemical potentials in the phase diagram of nuclear matter, within a modified and generalized hadron gas model, in order to consider phase transitions and to describe phenomena taking place outside the hadronic phase. We employ, in a phenomenological way, the Polyakov loop and scalar quark condensate order parameters, mass/temperature-scaled partition functions and enforce flavor conservation. We propose that the resulting variation of the heavy quark chemical potentials can be directly related to the quark deconfinement and chiral phase transitions. Then, the chemical potential of the strange and charm quark can be considered as an experimentally accessible "order parameter", probing the phase diagram of QCD.

1 Introduction

One of the main problems in the study of the phase transitions, occurring on the level of strong interactions, is finding an unambiguous observable, which would act as an experimentally accessible "order parameter" [1]. All proposed QGP signatures (strangeness enhancement, J/ψ suppression, dileptons, resonance shift and broadening, etc.) have already been observed in heavy ion collisions, however, we have seen, that they also occur, to some extent, in p – p or p – A interactions where no QGP production is theoretically expected. The physical quantity needed should exhibit a uniform behavior within each phase, but should change when a critical point is reached in a phase transition. It has been earlier suggested [2-4] that the chemical potential of strange quarks may be the sought-for macroscopic and therefore measurable thermodynamic quantity. The case of [2+1] flavors was thoroughly studied and it was shown that the change in the sign of the strange quark chemical potential, from positive in the hadronic phase to negative in the deconfined phase, may indeed be a unique indication of the deconfinement phase transition. Here we will review the basic aspects of the model and present the [2+2] flavors version, which is a generalization of the model with the inclusion of c-quark and charm hadrons.

2 Hadronic phase

Assuming that the system has attained thermal and chemical equilibration of four quark flavors (u, d, s, c), the partition function for the hadronic gas is written in the Boltzmann approximation:

\[ \ln Z_{HG}(T, V, \lambda_s, \lambda_c) = \ln Z_{HG}^{ud} + \ln Z_{HG}^{\text{strange}} + \ln Z_{HG}^{\text{charm}} \] (1)
where
\[ \ln Z_{HG}^{ud} = Z_m + Z_n(\lambda_q^3 + \lambda_q^{-3}) \] (2)

\[ \ln Z_{HG}^\text{strange} = Z_K(\lambda_q \lambda_s^{-1} + \lambda_q^{-1} \lambda_s) + Z_Y(\lambda_q^2 \lambda_s + \lambda_q^{-2} \lambda_s^{-1}) + Z_\Xi(\lambda_q^2 \lambda_s + \lambda_q^{-2} \lambda_s^{-1}) + Z_\Omega(\lambda_s^3 + \lambda_s^{-3}) \] (3)

and
\[ \ln Z_{HG}^\text{charm} = Z_D(\lambda_c \lambda_q^{-1} + \lambda_q \lambda_c^{-1}) + Z_{D_s}(\lambda_c \lambda_q^{-1} + \lambda_q \lambda_c^{-1}) + Z_{\Lambda_c, \Sigma_c}(\lambda_c \lambda_q^2 + \lambda_q^2 \lambda_c) + 2Z_{\Xi_c}(\lambda_q \lambda_c \lambda_s + \lambda_q^{-1} \lambda_s^{-1} \lambda_c^{-1}) + 2Z_{\Omega_c}(\lambda_s^2 \lambda_c + \lambda_s^{-2} \lambda_c^{-1}) \] (4)

is the partition function for the non strange, strange and charm sectors, respectively. The charm sector also includes strange/charm mesons and baryons that lead to a coupling of the fugacities \( \lambda_c, \lambda_s \). For simplicity we have assumed isospin symmetry \( \lambda_u = \lambda_d = \lambda_q \), while the one particle Boltzmann partition function is given by:
\[ Z_k(V, T) = \frac{VT^3}{2\pi^2} \sum_j g_j \left( \frac{m_j}{T} \right)^2 K_2 \left( \frac{m_j}{T} \right) \] (5)

The summation in Eq. (5) runs over the resonances of each hadron species with mass \( m_j \), and the degeneracy factor \( g_j \) counts the spin and isospin degrees of freedom of the \( j \)-resonance. For the strange hadron sector, kaons with masses up to 2045 MeV/c^2, hyperons up to 2350 MeV/c^2 and cascades up to 2252 MeV/c^2 are included, as well as the \( \Omega^- \) states at 1672 MeV/c^2 and 2252 MeV/c^2. For the charm hadron sector, we include purely charm mesons \( D^+, D^-, D^0 \) and baryons \( (\Lambda_c, \Sigma_c) \) as well as strange-charm mesons \( (D_{c}^{+}) \) and baryons \( (\Xi_c, \Omega_c) \) which contain both heavy quark flavors. All known charm resonances are taken into account with masses up to 2.7 GeV/c^2. To derive the Equation of State (EOS) of the hadron gas phase we simultaneously impose flavor conservation,
\[ < N_s - N_{\Xi} > = \frac{T}{V} \frac{\partial}{\partial \mu_s} \ln Z_{HG}(T, V, \lambda_q, \lambda_s, \lambda_c) = 0 \] (6)
\[ < N_c - N_{\Omega} > = \frac{T}{V} \frac{\partial}{\partial \mu_c} \ln Z_{HG}(T, V, \lambda_q, \lambda_s, \lambda_c) = 0 \] (7)

which reduce to a set of coupled equations:
\[ Z_K(\lambda_q^{-1} \lambda_s - \lambda_q \lambda_s^{-1}) + Z_Y(\lambda_q^2 \lambda_s - \lambda_q^{-2} \lambda_s^{-1}) + 2Z_{\Xi}(\lambda_q \lambda_s^2 - \lambda_q^{-1} \lambda_s^{-2}) + 3Z_{\Omega}(\lambda_s^3 - \lambda_s^{-3}) + Z_{\Xi_c}(\lambda_q \lambda_s \lambda_c - \lambda_q^{-1} \lambda_s^{-1} \lambda_c^{-1}) + 2Z_{\Omega_c}(\lambda_s^2 \lambda_c - \lambda_s^{-2} \lambda_c^{-1}) = 0 \] (8)

\[ Z_D(\lambda_c \lambda_q^{-1} - \lambda_q \lambda_c^{-1}) + Z_{D_s}(\lambda_c \lambda_q^{-1} - \lambda_q \lambda_c^{-1}) + Z_{\Lambda_c, \Sigma_c}(\lambda_c \lambda_q^2 - \lambda_q \lambda_c^2) + Z_{\Xi_c}(\lambda_q \lambda_s \lambda_c - \lambda_q^{-1} \lambda_s^{-1} \lambda_c^{-1}) + Z_{\Omega_c}(\lambda_s^2 \lambda_c - \lambda_s^{-2} \lambda_c^{-1}) = 0 \] (9)

The above conditions, define the relation between all quark fugacities and temperature in the equilibrated primordial state. In the HG phase with finite net baryon number density, the chemical potentials \( \mu_q, \mu_s \) and \( \mu_c \) are coupled through
the production of strange and charm hadrons. Due to this coupling $\mu_s, \mu_c > 0$ in the hadronic domain. A more elegant formalism describing the HG phase is the Strangeness-including Statistical Bootstrap model (SSBM) [5,6]. It includes the hadronic interactions through the mass spectrum of all hadron species, in contrast to other ideal hadron gas formalisms. The SSBM is applicable only within the hadronic phase, defining the limits of this phase. In the 3-flavor case, the hadronic boundary is given by the projection on the 2-dimensional $(T, \mu_q)$ phase diagram of the intersection of the 3-dimensional bootstrap surface with the strangeness-neutrality surface ($\mu_s = 0$). Note that the vanishing of $\mu_s$ on the HG borderline does not apriori suggest that $\mu_s = 0$ everywhere beyond. It only states that the condition $\mu_s = 0$ characterizes the end of the hadronic phase. Figure 1 exhibits the hadronic boundary for two heavy quark flavors, obtained by imposing the conditions $\mu_s = 0$ and $\mu_c = 0$ to Eq’s. (8), (9). Observe, that there exists an intersection point, at $T_{int} \sim 130$ MeV and $\mu_{q_{int}}^{T} \sim 325$ MeV. For an equilibrated primordial state (EPS) above this temperature, i.e $T > 130$ MeV, and low $\mu_q$ values, we observe that as the temperature decreases, the condition $\mu_c = 0$ is realized before the vanishing of $\mu_s$ (case I), whereas for $T < 130$ MeV and high $\mu_q$, the opposite effect takes place (case II). This behavior, may be of some importance towards a possible experimental identification of a color superconducting phase, which is realized at a low temperature and high density region of the phase diagram (case II).

Figure 1. The critical curves $\mu_s = 0$ and $\mu_c = 0$. We distinguish two cases depending on the location of the equilibrated primordial state (EPS).
3 Chirally symmetric QGP phase

The partition function for a four flavor Quark Gluon plasma has the form,
\[
\ln Z_{QGP}(T,V,\mu_{q,s,c}) = \frac{V}{T}\left[\frac{37}{90}\pi^2 T^4 + \frac{\mu_q^2 T^2}{2\pi^2} + \sum_{i=s,c} g \frac{m_i^0 T^2}{2\pi^2} (\lambda_i + \lambda_i^{-1}) K_2\left(\frac{m_i^0}{T}\right)\right]
\] (10)

where \(m^0_s, m^0_c\) is the current strange and charm quark masses respectively. Flavor conservation within the QGP phase yields \(\lambda_s = \lambda_c = 1\) or
\[
\mu_{QGP}^s(T,\mu_q,\mu_c) = \mu_{QGP}^c(T,\mu_q,\mu_s) = 0 \quad (11)
\]
throughout this phase. Here the two order parameters, the Polyakov loop \(<L>\) and the scalar quark density \(<\bar{\psi}\psi>\), have reached their asymptotic values. Note that the chirally symmetric quark-gluon plasma phase always corresponds to a vanishing heavy quark chemical potential.

4 Deconfined Quark Matter phase of \([2+2]\) flavors

We argue that, beyond the hadronic phase, an intermediate domain of deconfined yet massive and correlated quarks arises, according to the following qualitative picture: The thermally and chemically equilibrated primordial state at finite baryon number density, consists of the deconfined valence quarks of the participant nucleons, as well as of \(q-\bar{q}\) pairs, created by quark and gluon interactions. Beyond but near the HG boundary, \(T \geq T_d\), the correlation-interaction between \(q-\bar{q}\) is near maximum, \(\alpha_s(T) \leq 1\), a prelude to confinement into hadrons upon hadronization. With increasing temperature, the correlation of the deconfined quarks gradually weakens, \(\alpha_s(T) \rightarrow 0\), as color mobility increases. The mass of all (anti)quarks depends on the temperature and scales according to a prescribed way. The initially constituent mass decreases with increasing \(T > T_d\), and as the DQM region goes asymptotically into the chirally symmetric QGP phase, as \(T \rightarrow T_{\chi}\), quarks attain current mass. Thus, we expect the equation of state in the intermediate DQM region to lead to the EoS of the hadronic phase, Eq. (1), at \(T \leq T_d\), and to the EoS of the QGP, Eq. (6), at \(T \sim T_{\chi}\). In order to construct an empirical partition function for the description of the DQM phase, we use (a) the Polyakov loop \(<L> \sim e^{-F_q/T} \equiv R_d(T \geq T_d) = 0\) \(\rightarrow 1\) as \(T=T_d \rightarrow T_{\chi}\) and (b) the scalar density \(<\bar{\psi}\psi> \sim R_{\chi}(T \geq T_d) = 1\) \(\rightarrow 0\) as \(T=T_d \rightarrow T_{\chi}\). The first describes the quark deconfinement while the latter is associated with the quark mass scaling. We assume that above the deconfinement temperature, quarks retain some degree of correlation and can be considered as hadron-like states. Therefore, near \(T_d\) a hadronic formalism may still be applicable. This correlation/interaction gradually weakens, as a result of the progressive increase of color mobility. Each quark mass scales, decreasing from the constituent value to the current one as we reach the chiral symmetry restoration temperature \((T \rightarrow T_{\chi})\). Thus, we consider a temperature dependent mass for each quark flavor, approximated by:
\[
m^*_f(T) = R_{\chi}(T)(m_f - m^0_f) + m^0_f \quad (12)
\]
where $m_f$ and $m_0$ are the constituent and current quark masses respectively (the values $m_0^u = 5 MeV, m_0^d = 9 MeV, m_0^s = 170 MeV, m_0^c = 1.1 GeV$ have been used). In the same spirit, we approximate the effective hadron-like mass:

$$m_i(T) = R_\chi(T)(m_i - m_0^i) + m_0^i$$

(13)

where $m_i$ is the mass of each hadron in the hadronic phase and $m_0^i$ is equal to the sum of the hadron’s quarks current mass (for example $m_0^u = 175 MeV, m_0^s = 350 MeV$). In the partition function of the DQM phase, the former scaling is employed through the mass-scaled QGP partition function $\ln Z_{QGP}$, where all quark mass terms are given by Eq.(12), while the latter is used in the mass-scaled hadronic partition function $\ln Z_{HG}$, where all hadron mass terms are given by Eq.(12). Employing the described dynamics, we construct an empirical partition function for the DQM phase,

$$\ln Z_{DQM}(V,T,\{\lambda_f\}) = [1 - R_d(T)]\ln Z_{HG}(V,T,\{\lambda_f\}) + R_d(T)\ln Z_{QGP}^*(V,T,\{\lambda_f\})$$

for $f = q, s, c$.

The factor $[1 - R_d(T)]$ describes the weakening of the interaction of the deconfined quarks, while the factor $R_d(T)$ can be associated with the increase of color mobility as we approach the chirally symmetric QGP phase. The DQM partition function is a linear combination of the HG and QGP mass-scaled partition functions together with the general demand to describe both confinement and chiral symmetry restoration asymptotically. Note that below the deconfinement critical point $T < T_d$, $R_d(T) = 0$, leading to $\ln Z_{DQM} = \ln Z_{HG}$ (with constituent quarks), whereas at the chiral symmetry restoration temperature $T \sim T_s$, $R_d(T) = 1$ and $\ln Z_{DQM} = \ln Z_{QGP}$ (with current quark masses). In order to acquire the EoS of the DQM phase, we impose again the strangeness and charm neutrality conditions, leading to the set of equations respectively,

$$[1 - R_d(T)] [Z_{K^*}^f(\lambda_s \lambda_q^{-1} - \lambda_q \lambda_s^{-1}) + Z_{Y^*}^f(\lambda_s \lambda_q^2 - \lambda_q \lambda_s^2)] + 2Z_{\Xi^*}^f(\lambda_s \lambda_q - \lambda_q \lambda_s) + 3Z_{\Omega^*}^f(\lambda_q \lambda_c \lambda_q - \lambda_q \lambda_c \lambda_q^{-1} - \lambda_q \lambda_c \lambda_q^{-1}) + 2Z_{\Omega^*}^c(\lambda_q \lambda_c - \lambda_q \lambda_c) + R_d(T)g_s m_s^2 K_2 \left(\frac{m_s^2}{T}\right) (\lambda_s - \lambda_s^{-1}) = 0$$

(14)

and

$$[1 - R_d(T)] [Z_{D_s}^f(\lambda_c \lambda_q^{-1} - \lambda_q \lambda_c^{-1}) + Z_{D_c}^f(\lambda_c \lambda_s^{-1} - \lambda_s^{-1} \lambda_c)] + Z_{\Lambda_s}^f(\lambda_c \lambda_q \lambda_q^{-1} - \lambda_q \lambda_c \lambda_q^{-1} - \lambda_q \lambda_c \lambda_q^{-1}) + Z_{\Omega_c}^f(\lambda_c \lambda_c - \lambda_c \lambda_c) + R_d(T)g_c m_c^2 K_2 \left(\frac{m_c^2}{T}\right) (\lambda_c - \lambda_c^{-1}) = 0$$

(15)

which must be solved simultaneously. Note that because of the strange/charm hadrons $D_s, \Xi_c, \Omega_c$ there exists a coupling between the heavy quark fugacities $\lambda_s, \lambda_c$. By solving the above equations, for a given chemical potential $\mu_q$, we derive the variation of the strange and charm quark chemical potentials with temperature in the phase diagram.
5 Results for finite chemical potential

In the case of 3-flavors and finite density, we had neglected all terms involving c-quarks ($\lambda_c = 1$). In this case, only the variation of $\mu_s$ was considered and Figure 2 was derived. We observe that the strange quark chemical potential attains positive values in the hadronic phase, becomes zero upon deconfinement, it grows strongly negative in the DQM domain and finally returns to zero as the QGP phase is approached. It is important that $\mu_s$ behaves differently in each phase, as this is what we are looking for from the beginning in the search for an experimentally accessible "order parameter". The change in the sign of $\mu_s$ from positive in the hadronic phase to negative in the deconfined is an unambiguous indication of the quark deconfinement phase transition, as it is independent of assumptions regarding interaction mechanisms. In the case of [2+2] flavors the situation is slightly modified. Figure 3 exhibits the variation of the two correlated heavy quark chemical potentials with the temperature of the primordial state, as given by Eq's (14), (15). We observe that both are initially positive and then grow negative, although the change in their sign is realized at different temperatures, for example $\mu_s = 0$ at $T_d^s \sim 190 MeV$, while $\mu_c = 0$ at $T_d^c \sim 215 MeV$ for a fixed value of the fugacity $\lambda_q = 0.48$. However, this difference can be easily understood if we consider Figure 3 in the framework of Figure 1. As already discussed in Sec. 2, for an equilibrated primordial state (EPS) with $T > T_{int}$ and sufficiently low $\mu_q$, $\mu_c$ becomes zero earlier than $\mu_s$, as the system approaches hadronization (see Figure 1). This is the reason why $T_d^c > T_d^s$ in

Figure 2. Variation of $\mu_s$ with the temperature in the case of [2+1] quark flavors and different approximations or parameterizations of the order parameter $R_d(T)$.

Figure 3. For sufficiently high $\mu_q$ values and low temperatures, the opposite effect is present, i.e $\mu_c$ changes its sign at a lower temperature than the strange quark chemical potential. The magnitude of the difference $|T_d^c - T_d^s|$, will depend on the
exact location of the state in the phase diagram. The fact that the $\mu_s, \mu_c$ vanish at different temperatures, at the end of the respective hadronic domain, has further consequences, as it implies that there exists a quark "deconfinement region" rather than a certain critical line.

Figure 3. Plot of the strange and charm quark chemical potentials in the phase diagram for $\lambda_q = 0.48$. Notice that the change in their sign is realized at different temperatures.

6 Experimental data

Over the last years, data from several nucleus-nucleus collisions have been analyzed within thermal statistical models, employing the canonical and grand-canonical formalisms [7-11]. Table 1 summarizes some of the results for the quantities $T$, $\mu_q$ and $\mu_s$, which have been deduced after performing a fit to the experimental data. Figure 4 shows the phase diagram with the Ideal Hadron Gas (IHG) and SSBM $\mu_s = 0$ lines, as well as the location of the mean ($T, \mu_q$) values obtained for every collision. We observe that all interactions studied, are consistently situated inside the hadronic phase, defined by the IHG model and exhibit positive $\mu_s$. The sulfur-induced interactions, however are situated slightly beyond the hadronic phase defined by the SSBM. IHG calculations exhibit deviations from the SSBM as we approach the critical deconfinement point $T = T_d \sim 175$ MeV, where the S-S and S-Ag interactions are roughly located. Within the IHG model the condition $\mu_s = 0$ is satisfied at a higher temperature ($T \sim 200$ MeV), extending the hadronic phase to a larger region as can be seen in Figure 4. As a consequence, $\mu_s$ changes sign at a higher temperature also and this is the reason why $\mu_s > 0$ in the analysis of [11], although a temperature above deconfinement (according to the SSBM) is obtained. Therefore, an adjustment of the IHG curve to the SSBM boundary and a new fit to the data are necessary [12]. The data from RHIC at $\sqrt{s}=130, 200$ AGeV are not included in our discussion, since at such high energies $\mu_q$ is very small and $\mu_s \sim 0$ throughout the phase diagram. The observation of negative heavy quark chemical
potential requires a finite baryon density system.

![Diagram showing (T, µq) values of several interactions and their location in the phase diagram. The lines correspond to the hadronic boundary within the SSB and IHG models.]

Figure 4. (T, µq) values of several interactions and their location in the phase diagram. The lines correspond to the hadronic boundary within the SSB and IHG models.

7 Conclusions

On the basis of the present analysis, we conclude that the heavy quark chemical potentials behave differently in each region (HG-DQM-QGP) of the phase diagram and, therefore, they can serve as a probe of the phase transitions. This is the first proposal of such an experimentally accessible "order parameter" that holds for a finite baryon density state. The appearance of negative values of µs and µc, is a well-defined indication of the quark deconfinement phase transition, at T=Tc, which is free of ambiguities related to microscopic effects of the interactions. It is important to add, that the observation of negative heavy quark chemical potentials would be also a clear evidence for the existence of the proposed DQM phase, meaning that chiral symmetry and deconfinement are apart at finite density. Until now, there is no known argument from QCD that the two transitions actually occur at the same temperature. Au+Au collisions at intermediate energies, for example 30 ≤ √s ≤ 90 AGeV, should be performed to experimentally test our proposals.

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Table 1. Deduced values for $T$, $\mu_q$, $\mu_s$ from several thermal models and fits to experimental data for several interactions.

| Interaction/Experiment                      | $T$ (MeV)  | $\mu_q$ (MeV) | $\mu_s$ (MeV) |
|---------------------------------------------|------------|---------------|---------------|
| Si+Au(14.6 AGeV)/E802                       |            |               |               |
| Reference[4]                                | 134±6      | 176±12        | 66±10         |
| Reference[9]                                | 135±4      | 194±11        | 66±10         |
| Mean                                        | 135±3      | 182±5         |               |
| Pb+Pb(158 AGeV)/NA49                        |            |               |               |
| Reference[4]                                | 146±9      | 74±6          | 22±3          |
| Reference[9]                                | 158±3      | 79±4          | 25±4          |
| Reference[7]                                | 157±4      | 81±7          | 23±2          |
| Mean                                        | 157±3      | 78±3          |               |
| Pb+Pb(40 AGeV)/NA49                         |            |               |               |
| Reference[4]                                | 147±3      | 136±4         | 35±4          |
| Reference[*]                                | 150±8      | 132±7         |               |
| Mean                                        | 149±9      | 134±8         |               |
| S+S(200 AGeV)/NA35                          |            |               |               |
| Reference[10]                               | 182±9      | 75±6          | 14±4          |
| Reference[11]                               | 181±11     | 73±7          | 17±6          |
| Reference[8]                                | 202±13     | 87±7          | 16±7          |
| Mean                                        | 188±6      | 78±4          |               |
| S+Ag(200 AGeV)/NA35                         |            |               |               |
| Reference[10]                               | 180±3      | 79±4          | 14±4          |
| Reference[11]                               | 179±8      | 81±6          | 16±5          |
| Reference[8]                                | 185±8      | 81±7          | 16±8          |
| Mean                                        | 181±4      | 80±3          |               |

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