Resonance States in an effective Chiral Hadronic Model

P. Rau¹,²,*, J. Steinheimer², S. Schramm¹,², H. Stöcker¹,³

1 Institut für Theoretische Physik, Goethe-Universität, Max-von-Laue-Str. 1, 60438 Frankfurt am Main, Germany
2 Frankfurt Institute for Advanced Studies (FIAS), Ruth-Moufang-Str. 1, 60438 Frankfurt am Main, Germany
3 GSI Helmholtzzentrum für Schwerionenforschung GmbH, Planckstr. 1, 64291 Darmstadt, Germany

Abstract: With an effective chiral flavour SU(3) model we show the effect of hadronic resonances on the QCD phase diagram. We state that varying the resonance couplings to the scalar and vector fields affects the order and location of the phase transition, the possible existence of a critical end point (CEP), and the thermodynamic properties. We present (strange) quark number susceptibilities at zero baryochemical potential and at three different points at the phase transition. Comparing results to lattice QCD, we state that reasonable large vector couplings limit the phase transition to a smooth crossover ruling out a CEP.

PACS (2008): 12.38.-t, 11.30.Rd, 14.20.Gk

Keywords: Effective Model • Hadronic resonances • QCD phase diagram • Susceptibilities

© Versita Warsaw and Springer-Verlag Berlin Heidelberg.

1. Introduction

Both on the theoretical and on the experimental side there are great efforts to study the largely unknown properties of the QCD phase diagram with special attention to the phase transitions of strongly interacting matter. Most interesting are the transitions that restore chiral symmetry and that generate deconfinement from hadrons to a Quark Gluon Plasma (QGP) at high temperatures and densities. Experimentally different regions of the phase diagram are studied in ultra-relativistic heavy ion collisions at different beam energies. Since QCD cannot be solved analytically, there are mainly two groups of theoretical approaches to QCD. On the one hand there are lattice QCD models, which are mostly limited to vanishing baryochemical potentials, and on the other there are effective models describing specific properties of QCD matter. Here, we present an effective chiral flavour SU(3) model with hadronic degrees of freedom including heavy resonance states and study the properties of the phase diagram compared to lattice QCD results.

* E-mail: rau@th.physik.uni-frankfurt.de
2. Model

Our SU(3)-flavour $\sigma$-$\omega$-model with a non-linear realization of chiral symmetry [1] is based on the mean field Lagrangian $\mathcal{L} = \mathcal{L}_{\text{kin}} + \mathcal{L}_{\text{int}} + \mathcal{L}_{\text{mes}}$. With the hadrons’ kinetic energy in the first term, the interaction of baryons with scalar $\sigma$, $\zeta$ and vector mesons $\omega$, $\phi$ is described by

$$\mathcal{L}_{\text{int}} = -\sum_i \bar{\psi}_i \left( m^*_i + g_\omega \gamma_0 \omega^0 + g_\phi \gamma_0 \phi^0 \right) \psi_i.$$  \hspace{1cm} (1)

The index $i$ includes the baryon octet, decuplet, and heavier resonances with masses $m \leq 2.6$ GeV with a minimum three star rating in the Particle Data Book. The term $\mathcal{L}_{\text{mes}} = \mathcal{L}_{\text{vec}} + \mathcal{L}_0 + \mathcal{L}_{\text{ESB}}$ includes the vector and the scalar mesons’ self interactions and the explicit chiral symmetry breaking. The effective baryon masses $m^*_i = g_\omega \sigma + g_\zeta \zeta + \delta m_i$ are generated by the coupling of baryons to the scalar meson fields $\sigma$, $\zeta$ and by an explicit mass $\delta m_i$. At high temperatures and baryonic densities this formalism leads to smaller baryon masses and thereby to the restoration of chiral symmetry due to the decreasing $\sigma$-field. The couplings of the baryon octet to the mesonic fields and the mesonic potential are fixed to reproduce the tabulated vacuum masses and well-known nuclear ground state properties. The coupling strengths of the baryonic resonances are set proportional to the nucleon couplings $g_N$ via the parameters $r_{s,v}$ according to $g_B \sigma, \omega = r_{s,v} \cdot g_N \sigma, \omega$ and $g_B \zeta, \phi = r_{s,v} \cdot g_N \zeta, \phi$.

For simplicity and for generating the effective baryon masses dynamically by the scalar fields, we keep the scalar coupling fixed at $r_s = 0.97$. Thereby, we are able to reproduce the particles’ vacuum masses, except for the explicit mass $\delta m_i$, and to ensure a smooth crossover transition at zero baryochemical potential $\mu_B = 0$. However, the vector coupling parameter $r_v$, which controls the abundances of the baryonic resonances at $\mu_B \neq 0$, is varied in order to study its impact on the calculated phase diagram and the thermodynamic properties. In contrast to this interacting hadron resonance gas (HRG), we also perform calculations for an ideal non-interacting HRG. In this particular case $r_{s,v}$ are set to zero and the masses of all particles are fixed at their tabulated vacuum expectation value.
The grand canonical potential $\Omega/V = -\mathcal{L}_\text{int} - \mathcal{L}_\text{meson} + \Omega_\text{th}$ includes the thermal contribution of all hadrons (i.e. all particles and antiparticles) in the model. The effective baryochemical potential is defined as $\mu^*_i = \mu_i - g_i^\omega \omega - g_i^\phi \phi$.

From the grand canonical potential all thermodynamic quantities follow, i.e. the pressure $p$, the energy and entropy densities $e$, $s$, and the particle densities $\rho_i$.

### 3. Results

The normalised chiral order parameter $\sigma/\sigma_0(T)$ from our model is shown in Fig. 1 (left) together with lattice data from different actions and lattice spacings [2]. With $r_v$ fixed, $\sigma/\sigma_0$ of the interacting HRG (blue dotted line) exhibits a smooth crossover and a critical temperature $T_c = 164$ MeV being in line with recent lattice data for the continuum extrapolated HISQ and stout actions. Analysing $\sigma/\sigma_0$ in the whole $T - \mu$-plane we get the phase transition lines shown in Fig. 1 (right) for different values of $r_v$. Here, dashed lines depict crossover transitions, dots a CEP, and full lines first order phase transitions. We find that for reasonable values of $r_v \geq 0.4$, for which the nuclear ground state is correctly located in the chirally broken phase, the CEP moves to low temperatures or vanishes completely ($r_v > 0.6$) in favour of a broad crossover transition.

Figure 2 shows the interaction measure $(e - 3p)/T^4$ (left) and the strange susceptibility $\chi_s/T^2$ (right) reflecting strange quark fluctuations [3] both as functions of $T$. As expected, compared to the non-interacting HRG the interacting HRG (blue lines) shows a much faster increase of both quantities at $T_c$ due to more degrees of freedom. Both quantities from the interacting HRG are in qualitatively good agreement with the stout continuum limit.

Studying the second and fourth order quark number susceptibility coefficients $c_2$, $c_4$ at $\mu_B = 0$ (Fig. 3 (a) and (b)), we find that fluctuations at $T_c$ get massively suppressed by large couplings to the repulsive vector field $r_v$. This is also reflected in the susceptibility ratio $c_4/c_2$ (Fig. 3 (c)) which only comes close to lattice data for reasonably high values of $r_v \geq 0.8$. This again supports our previously mentioned preference for larger values $r_v$.
Figure 3. (Color online) Susceptibility coefficients \( c_2 \), \( c_4 \) ((a) and (b)) and the ratio \( c_4/c_2 \) as a function of \( T/T_c \) at \( \mu_B = 0 \) (c) for the interacting (blue lines) and non-interacting (black lines) HRG for \( r_v = 0.8 \) and at different points (CEP - \( \mu_B = 216 \) MeV; CO - \( \mu_B = 30 \) MeV; PT - \( \mu_B = 490 \) MeV) on the phase transition for \( r_v = 0.9 \) with lattice data (d).

and thus limits the phase transition to a broad crossover in the whole \( T - \mu \)-range (cf. Fig. 1 (right)). For the quark number susceptibilities we think, a comparison to not yet available lattice data with continuum extrapolated stout action would be very promising. A strong suppression of the susceptibility coefficients is also seen at \( \mu_B \neq 0 \). When calculating the ratio \( c_4/c_2 \) crossing the phase transition at different values of \( \mu_B \) (Fig. 3 (d)), we find that fluctuations are much smaller compared to \( \mu_B = 0 \) which is again due to the repulsive vector interactions.

References

[1] J. Boguta and H. Stöcker, Phys.Lett. B120, 289 (1983); P. Papazoglou et al., Phys.Rev. C57, 2576 (1998); Phys.Rev. C59, 411 (1999); P. Rau et al., (2011), arXiv:1109.3621.

[2] A. Bazavov et al., Phys. Rev. D80, 014504 (2009); A. Bazavov and P. Petreczky, J. Phys. Conf. Ser. 230, 012014 (2010); (2010), arXiv:1009.4914; M. Cheng et al., Phys. Rev. D79, 074505 (2009); Phys. Rev. D81, 054504 (2010); S. Borsanyi et al., JHEP 1009, 073 (2010); Y. Aoki et al., JHEP 0906, 088 (2009); S. Ejiri et al., Prog. Theor. Phys. Suppl. 153, 118 (2004); G. Endrodi et al., JHEP 1104, 001 (2011).

[3] V. Koch, in Landolt-Börnstein, Relativistic Heavy Ion Physics, Vol. I/23, edited by R. Stock (Springer, Berlin,2010), arXiv:0810.2520.