Physics Letters B 778 (2018) 419–425

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Holographic QCD phase diagram with critical point from Einstein–Maxwell-dilaton dynamics

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ARTICLE INFO

Article history:
Received 24 February 2017
Received in revised form 27 October 2017
Accepted 18 January 2018
Available online 3 February 2018
Editor: J.-P. Blaizot

Keywords:
Gravity dual
Holography
Quark–gluon plasma
Critical point

ABSTRACT

Supplementing the holographic Einstein–Maxwell-dilaton model of [1,2] by input of lattice QCD data for $2+1$ flavors and physical quark masses for the equation of state and quark number susceptibility at zero baryo-chemical potential we explore the resulting phase diagram over the temperature-chemical potential plane. A first-order phase transition sets in at a temperature of about 112 MeV and a baryo-chemical potential of 612 MeV. We estimate the accuracy of the critical point position in the order of approximately 5–8% by considering parameter variations and different low-temperature asymptotics for the second-order quark number susceptibility. The critical pressure as a function of the temperature has a positive slope, i.e. the entropy per baryon jumps up when crossing the phase border line from larger values of temperature/baryo-chemical potential, thus classifying the phase transition as a gas–liquid one. The updated holographic model exhibits in- and outgoing isentropes in the vicinity of the first-order phase transition.

1. Introduction

The QCD phase diagram exhibits potentially a large variety of structures [3–6]. Either originating from extrapolations of weak-coupling results or being suggested by models (most notably Nambu–Jona–Lasinio (cf. [7]), linear sigma/quark-meson [8] models in numerous variants), various phases of strongly interacting matter may occur, such as color superconductors (cf. [9–11]), or quarkyonic matter (cf. [12]), or chirally restored phases (cf. [13]), or color-flavor locked structures (cf. [14]).

While the gas–liquid (GL) first-order phase transition (FOPT) in nuclear matter seems to be well established since some time [15–20], the hadron–quark (HQ) deconfinement transition still offers a few challenges. At very small or zero net-baryon density corresponding to a small chemical potential (μ)-to-temperature (T) ratio, $\mu/T \ll 1$, the HQ transition is established as a crossover in $2+1$ flavor lattice QCD with physical quark masses [21,22] at a characteristic scale of $T_c = \mathcal{O}(150\text{ MeV})$. The popular Columbia plot [23] sketches qualitatively the options of the phase structure in dependence of the u, d, s quark masses $m_{u,d,s}$. For instance, in the chiral limit, $m_{u,d,s} \to 0$, or the opposite infinitely heavy quark-mass limit, $m_{u,d,s} \to \infty$, the deconfinement transition is a FOPT. Due to the sign problem of the fermionic determinant the ab initio lattice QCD evaluations are not yet conclusive with respect to the confinement and chiral restoration transition(s) at non-zero baryo-chemical potential, in particular for $\mu/T > 2$. Some methods try to avoid or circumvent the sign problem (cf. [24]), e.g. by evaluations at imaginary μ (which need a prescription of $i\mu \to \mu$) or a Taylor expansion in powers of μ/T with coefficients calculated at $\mu = 0$ (which needs statements on the convergence [25]), or the reweighting method (which needs statements on the density and parameter ranges to incorporate the sign and overlap problem [26]). Other approaches are based on the complex Langevin method [27,28] (see [29] for recent developments) or a recent proposal for a path optimization method [30], which is based on the Lefschetz-thimble path-integral method [31].

The pertinent uncertainties make the region of larger μ/T interesting. A particularly interesting option is the possibility of a (critical) end point (CEP) of a curve of FOPTs, e.g. $T_c(\mu)$, setting in at (T_{CEP}, μ_{CEP}) and running toward the $T = 0$ axis when imaging the phase diagram in the T–μ plane.

The CEP coordinates are yet fairly unconstrained. Plugging model results and QCD-related extrapolations together one arrives at some less conclusive scatter plot (cf. e.g. [24]). Advanced lattice QCD approaches disfavor a CEP position at $T/T_c(\mu = 0) > 0.9$ and $\mu/T \le 2$ [25].

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https://doi.org/10.1016/j.physletb.2018.01.053

Experimentally, there are dedicated programs aiming at pinning down the CEP location. For instance, the beam energy scan at RHIC [32] gave hints on some features in the beam energy dependence of selected observables which have been interpreted as CEP signature (cf. [33]). In [34] another view has been launched with the conclusion of having also seen CEP indications. Furthermore, the SHINE (NA61) collaboration at CERN-SPS is also systematically seeking CEP effects [35]. Experiments planned at FAIR and NICA and J-PARC [36] are analogously driven by CEP searches, analogously as goals by the CBM collaboration [37,38], and the MPD group [39].

Given that challenges from both theory and experiment one can ask whether further theoretical model classes beyond the above mentioned approaches could be useful in exploring the hypothetical FOPT emerging from a CEP. Holographic models, advancing the seminal AdS/CFT correspondence [40–42], are thought to mimic essential QCD properties in the strong-coupling regime [43–47] and thus may serve as suitable candidates for such an enterprise. Holographic bottom-up approaches coupled to a self-interacting dilaton with nontrivial potential were particularly successful to describe nonconformal properties of the quark–gluon plasma and QCD [48–51]. In [1,2] a model formulation has been put forward which displays a critical point in the T–μ plane. While [1,2] focuses on CEP properties and an outline of some transport coefficients, [52,53] employed that holographic model to investigate thermodynamics and further transport quantities at small μ/T, however, the question of the CEP position, based on an adjustment to recent lattice data, and properties of phase diagrams were not addressed. The model rests on the coupled Einstein–Maxwell-dilaton (EMd) dynamics and can be adjusted to QCD thermodynamics, i.e. the equation of state (EoS) and quark number susceptibility at $\mu = 0$. The resulting phase structure is the topic of our present paper. We feel that an update of [1,2] is timely since by now consistent and more precise lattice QCD data are at our disposal. In fact, we find some qualitatively important modifications in comparison to [1,2] w.r.t. the pattern of isentropes in the phase diagrams as well as the position of the CEP.

With respect to the discussion in [54], a FOPT curve is specified by further peculiarities: it can be related either to a GL type or to a HQ type transition. For a discussion contrasting features of GL and HQ phase transitions we refer the interested reader to [54–57], where the notions of entropic vs. enthalpic transitions as well as congruent and non-congruent material changes are exemplified and representations in other variables than T–μ are exhibited. Such different FOPTs can matter significantly in core-collapse supernova explosions as discussed in some detail in [58]. Motivated by such a relation to astrophysical aspects of the phase structure of strongly interacting matter – not only touching core-collapse dynamics but also neutron (quark core) stars – we unravel here the phase structure of the holographic EMd model. It turns out that the EMd model with adjustments to QCD input belongs to the GL class. That is across the phase boundary both the baryon density n and the entropy density s jump when considering the stable phases. For the GL transition, the entropy per baryon s/n drops down when going into μ or T direction, while at the HQ transition s/n jumps up, according to expectations in [54]. According to the Clausius–Clapeyron equation one finds the critical pressure $p(T, \mu_c(T))$ either with positive slope (GL transition) or with negative slope (HQ transition).[1]

[1] Obviously, the resulting behavior of the pressure at the FOPT at smaller temperatures is markedly depending on these details, with impact on the stiffness of the EoS which in turn governs the possibility of a third family of compact stars or twin configurations [59–62], on which the options for core-collapse supernova explosions according to [58] (and further references therein) depend on.

Our paper is organized as follows. In section 2 we recall the holographic EMd model. The numerical adjustment to lattice QCD data at $\mu = 0$ is described in section 3 and the numerical results for the phase diagrams are presented in section 4, including an analysis of the impact of different assumptions for the susceptibility at small temperatures. We summarize in section 5.

2. Recalling the holographic EMd model

The holographic model of gravity of a 5-dimensional Riemann space sourced by the coupled Maxwell-dilaton fields is defined in [1,2] by the action

$$S = \frac{1}{2\kappa_5^2} \int d^5x \sqrt{-g} \left(R - \frac{1}{2}\partial^\mu \phi \partial_\mu \phi - V(\phi) - \frac{f(\phi)}{4} F_{\mu\nu}^2 \right) + S_{GH}, \tag{1}$$

where R is the Einstein–Hilbert part, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ with $A_\mu dx^\mu = \Phi dt$ stands for the Abelian gauge field à la Maxwell, and ϕ is a real scalar (dilaton) with self-interaction described by the so called potential $V(\phi)$. The Maxwell field and dilaton are coupled by a dynamical strength function $f(\phi)$. The Gibbons–Hawking term S_{GH} for a consistent formulation of the variational problem is not needed explicitly in our context. The "Einstein constant" κ_5 is taken as a model parameter. The ansatz for the infinitesimal line element squared

$$ds^2 = e^{2A(r;r_H)}\left(-h(r;r_H)dt^2 + d\vec{x}^2\right) + \frac{e^{2B(r;r_H)}dr^2}{h(r;r_H)} \tag{2}$$

highlights that (i) only the dynamics in bulk direction r is considered and (ii) a horizon is admitted at $r = r_H$ by a simple zero of the blackness function h. By a gauge choice, one can achieve $B = 0$ and $r_H = 0$. We solve the field equations following from (1), (2) with the technique described in [1,2]. In a nutshell: One has to numerically integrate from $r_H + \epsilon$ towards the boundary at $r \to \infty$. Requiring regularity of A, h, ϕ, Φ at the horizon $r = r_H$, defined by $h(r_H; r_H) = 0$, series solutions for any these functions can be obtained, which yield the initial conditions for the integration. After fixing all gauge redundancies the two remaining independent quantities parametrizing the solutions are $\phi_0 \equiv \phi(r_H, r_H)$ and $\Phi_1 \equiv \frac{\partial \Phi}{\partial r}\big|_{r_H}$. It follows from the horizon expansion of A that Φ_1 is bounded, $\Phi_1 < \Phi_1^{max} \equiv \sqrt{-\frac{2V(\phi_0)}{f(\phi_0)}}$. Close to the boundary, the following expansions in powers of $e^{-\alpha(r)} \equiv \exp[-\frac{r}{L\sqrt{h_0^\infty}} - A_0^\infty]$ are valid: $h(r) = h_0^\infty + \ldots$, $A(r) = \alpha(r) + \ldots$, $\Phi(r) = \Phi_0^\infty + \Phi_2^\infty e^{-2\alpha(r)} + \ldots$, and $\phi(r) = \phi_A e^{-(4-\Delta)\alpha(r)} + \phi_B e^{-\Delta\alpha(r)} + \ldots$. The expansion of ϕ assumes $L^2 V(\phi) = -12 + \frac{1}{2}[\Delta(\Delta - 4)]\phi^2 + \ldots$ for $\phi \to 0$.[2] By the standard AdS/CFT dictionary, ϕ_A is the source and ϕ_B the expectation value of the boundary theory operator dual to ϕ. Then one obtains the thermodynamic quantities temperature T, entropy density s, baryo-chemical potential μ and baryon density n as

$$T = \lambda_T \frac{1}{4\pi \phi_A^{1/(4-\Delta)} \sqrt{h_0^\infty}}, \tag{3}$$

$$s = \lambda_s \frac{2\pi}{\phi_A^{3/(4-\Delta)}}, \tag{4}$$

[2] This means we are considering a relevant operator in the boundary theory with scaling dimension $\Delta < 4$; see [63] for a different choice of potential asymptotics that correspond to a marginal operator.