The QCD phase transition in a fully dynamical model of heavy-ion collisions

Christoph Herold[1,2], Marlene Nahrgang[2,3], Igor Mishustin[2,4], Marcus Bleicher[1,2]

[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] Department of Physics & CTMS, Duke University, Durham, NC 27708, USA
[4] Kurchatov Institute, National Research Center, 123182 Moscow, Russia

E-mail: herold@fias.uni-frankfurt.de

Abstract. Experimental signals for a possible QCD critical point and first-order phase transition are strongly influenced by the rapid nonequilibrium dynamics during a heavy-ion collision. In order to estimate and understand these effects we study the cooling through the phase transition within a nonequilibrium chiral fluid dynamics model. The order parameters for the chiral and deconfinement transition are explicitly propagated, taking into account dissipation and fluctuation stemming from the interaction with a quark-antiquark fluid. In studies of single events, we demonstrate how the formation of domains in net-baryon density at the first-order phase transition leads to a clear enhancement of higher flow harmonics. For the detection of the critical point it is crucial that the relevant signal survives the rapid dynamics. We observe critical slowing down and long-wavelength fluctuations in the vicinity of the critical point.

1. Introduction

Exploring the phase structure of strongly interacting matter is one of the primary goals of heavy-ion physics. The implementation of quantum chromodynamics (QCD) on the lattice predicts a crossover transition at small baryochemical potential $\mu_{\rm B}$ [1] with a pseudo-critical temperature in the range between $150 - 170$ MeV [2]. The concept of a critical point (CP) and a first-order phase transition at large values of $\mu_{\rm B}$ is mainly supported by mean-field effective models [3] and functional methods such as Dyson-Schwinger equations [4] or renormalization group techniques [5]. Experimentally, fluctuations are of crucial importance for the detection of any of these transition scenarios. At the CP, a diverging correlation length in equilibrium would lead to diverging event-by-event fluctuations of conserved quantities [6]. On the other hand, a different type of fluctuation is expected at a dynamical nonequilibrium first-order phase transition, where spinodal instabilities could foster spatial inhomogeneities within single events [7, 8]. We address these issues within the nonequilibrium chiral fluid dynamics model trying to find hints for both critical behavior and spinodal decomposition.

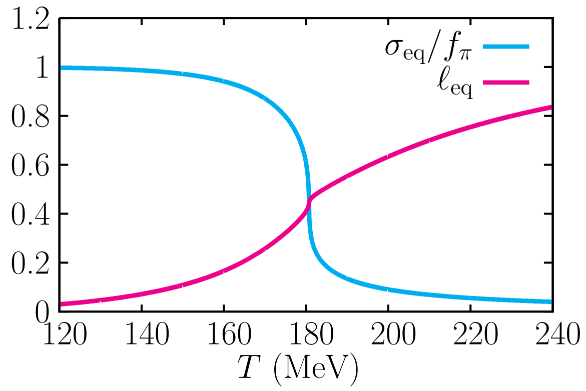

Figure 1. Equilibrium values of the sigma and Polyakov loop fields for a critical point scenario with $g = 3.52$.

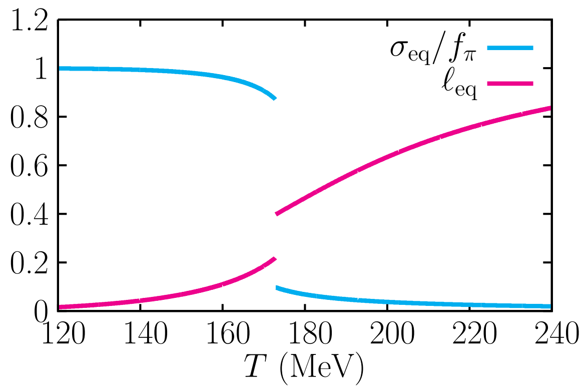

Figure 2. Equilibrium values of the sigma and Polyakov loop fields for a scenario with a first-order phase transition with $g = 4.7$.

2. Coupling fluid dynamics to chiral fields and Polyakov loop

Starting from the Polyakov-Quark-Meson model [9], we obtain the dynamics of the sigma field, the chiral order parameter, by adding damping and stochastic noise to the usual Euler-Lagrange equation

$$\partial_\mu \partial^\mu \sigma + \eta_\sigma(T)\partial_t \sigma + \frac{\partial V_{\rm eff}}{\partial \sigma} = \xi_\sigma \ . \tag{1}$$

This can be explicitly derived using the two-particle irreducible effective action approach [10], which gives a temperature-dependent damping coefficient $\eta_\sigma(T)$ vanishing only around the CP, and a dissipation-fluctuation relation. Here, $V_{\rm eff}$ is the grand canonical potential

$$V_{\rm eff} = U(\sigma) + \mathcal{U}(\ell) + \Omega_{\rm q\bar{q}} \ . \tag{2}$$

It is composed of the chiral potential U describing spontaneous breaking of chiral symmetry, the effective Polyakov loop potential $\mathcal{U}$ and the mean-field quark-antiquark contribution $\Omega_{q\bar{q}}$ from integrating out the quark degrees of freedom in the path integral formulation of the partition function.

A rigorous derivation of the equation of motion for the Polyakov loop is not possible. This quantity is only defined in Euclidean time and we do not understand the real-time equivalent. All we can do is make some reasonable phenomenological estimate. From the mean-field approximation we would have $\partial V_{\rm eff}/\partial \ell = 0$, which forces ℓ to be equal to its thermal expectation value at all times. To gain an independent dynamical description we augmented the equation of motion with damping and noise terms, giving [11]

$$\eta_\ell \partial_t \ell + \frac{\partial V_{\rm eff}}{\partial \ell} = \xi_\ell \ . \tag{3}$$

Thereby, we found our results to be independent of the choice of η_ℓ for a wide range of applications. We chose the value of $\eta_\ell = 5/{\rm fm}$. Both noise fields ξ_σ and ξ_ℓ are assumed to be Gaussian and white.

The quark fluid is described by an ideal stress-energy tensor $T^{\mu\nu}_{\rm q} = (e+p)u^\mu u^\nu - p g^{\mu\nu}$, where pressure and energy density are determined via $p = -\Omega_{\rm q\bar{q}}$ and $e = T\partial p/\partial T - p + \mu \partial p/\partial \mu$. The coupling to the order parameter fields σ and ℓ is achieved through source terms S^ν_σ and S^ν_ℓ,

$$\partial_\mu T^{\mu\nu}_{\rm q} = S^\nu_\sigma + S^\nu_\ell \ , \tag{4}$$

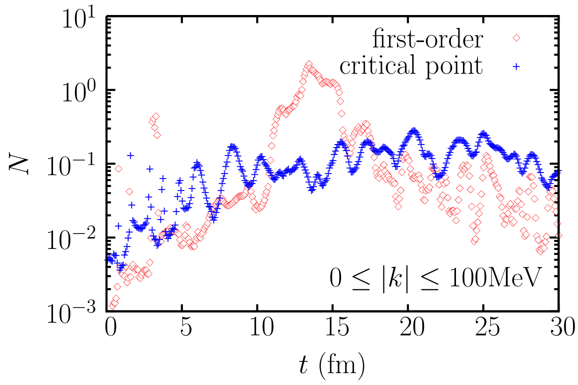

Figure 3. Fluctuation intensity of the sigma field for $0 \le |k| \le 100$ MeV.

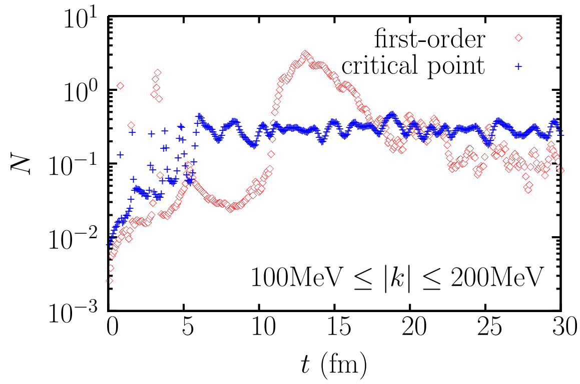

Figure 4. Fluctuation intensity of the sigma field for $100 \le |k| \le 200$ MeV.

which ensures conservation of the total energy of the coupled system [11, 12]. It is important to note here, that the local pressure of the fluid depends on the local values of the fields, so we do not use an equilibrium equation of state. Conservation of the total quark number is ensured by the continuity equation $\partial_\mu n^\mu = 0$.

3. Enhancement of low-momentum modes at the critical point

We investigate the dynamics of the order parameter fields after a temperature quench. For this purpose we put the system in a box of finite size with periodic boundary conditions. In this setup, pressure gradients are negligibly small and the dynamics is dominated by the Langevin dynamics of the fields. We restrict ourselves to the case of vanishing baryochemical potential and vary the strength of the transition by changing the quark-meson coupling g. In principle this quantity has to be chosen such that $g\sigma$ in vacuum reproduces the constituent quark mass, resulting in $g \sim 3.3$ and a crossover chiral and deconfinement transition. For a value of $g = 3.52$ we obtain a CP and for larger values of g the transition is of first-order type. Figs. 1 and 2 illustrate this through the equilibrium values $\sigma_{\rm eq}$ and $\ell_{\rm eq}$ as a function of temperature.

We let the system in the box relax in the vicinity of the respective transition temperature for the case of a CP and a first-order phase transition. During the evolution we focus on the behavior of sigma field fluctuations which we analyze through their intensity. This quantity is given by a combination of creation and annihilation operators which in equilibrium resembles the number of excited particles from the quantum field. It reads

$$\frac{\mathrm{d}N}{\mathrm{d}^3k} = \frac{a_k^\dagger a_k}{(2\pi)^3 2\omega_k} = \frac{\omega_k^2 |\delta\sigma_k|^2 + |\partial_t \sigma_k|^2}{(2\pi)^3 2\omega_k} \ , \tag{5}$$

with the Fourier transformed field $\delta\sigma_k$ and the time derivative $\partial_t\sigma_k$. The corresponding energy of the k-the mode is given by $\omega_k = \sqrt{m_\sigma^2 + \vec{k}^2}$. The intensities integrated over the low-momentum ranges from 0 to 100 MeV and from 100 to 200 MeV are plotted in Figs. 3 and 4, respectively. In both figures we can make out significant differences between the two types of transition: For the CP, the field needs less time to reach the equilibrium value, nevertheless, fluctuations remain strong here with a clear enhancement compared to the first-order phase transition towards the end of the evolution. On the other hand, we find large fluctuations when the system undergoes the first-order phase transition, induced by spinodal instabilities around the potential barrier separating the two competing minima. After equilibration, fluctuations are significantly weaker than near the CP.

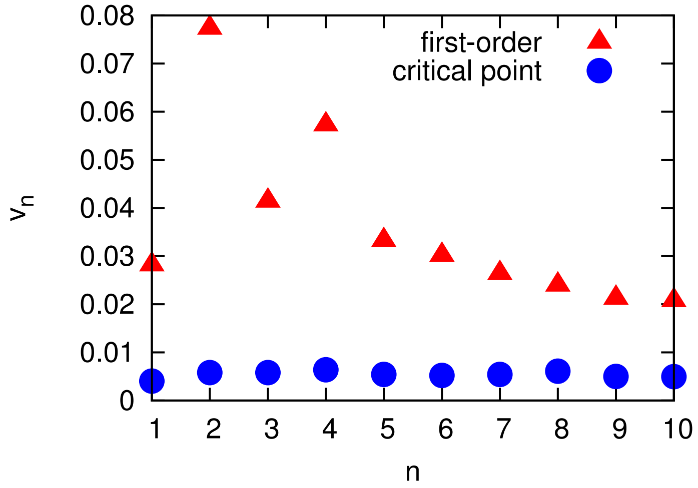

Figure 5. Event-averaged harmonic coefficients from the azimuthal net-baryon number distribution. We find a strongly anisotropic expansion for at first-order transition, while small coefficients indicate a homogeneous spherical shape at the critical point.

4. Anisotropic flow at large baryon densities

In order to study what happens during a heavy-ion collision, we consider a freely expanding blob of hot and dense quark matter cooling through the first-order phase transition at large $\mu_{\rm B}$. Here, we find that a metastable phase dynamically fragments into small droplets, leading to inhomogeneities in the azimuthal baryon number distribution within single events [13]. Averaged over many events, we find a clear enhancement of higher harmonics in comparison with an evolution through the CP, cf. Fig. 5. Note here that we used spherical initial conditions instead of a more realistic ellipsoidal shape to illustrate the mere influence of the phase transition.

In the future, we aim at investigating more realistic scenarios by introducing initial conditions from transport models and a hadronic freeze-out. We expect that the enhancement of the harmonic coefficients at the first-order phase transition leads to strongly enhanced flow.

Acknowledgements

The authors thank Stefan Leupold and Carsten Greiner for fruitful discussions and Dirk Rischke for providing the SHASTA code. This work was supported by the Hessian LOEWE initiative Helmholtz International Center for FAIR. I. M. acknowledges partial support from the grant NSM-215.2012.2 (Russia).

References

[1] Aoki Y, Endrodi G, Fodor Z, Katz S D and Szabo K K, *Nature* **443** (2006) 675
[2] Borsanyi S *et al.* [Wuppertal-Budapest Collaboration], *JHEP* **1009** (2010) 073
[3] Scavenius O, Mocsy A, Mishustin I and Rischke D H, *Phys. Rev. C* **64** (2001) 045202
[4] Fischer C S and Luecker J, *Phys. Lett. B* **718** (2013) 1036
[5] Herbst T K, Pawlowski J M and Schaefer B J, *Phys. Rev. D* **88** (2013) 014007
[6] Stephanov M A, Rajagopal K and Shuryak E V, *Phys. Rev. D* **60** (1999) 114028
[7] Scavenius O and Dumitru A, *Phys. Rev. Lett.* **83** (1999) 4697
[8] Steinheimer J and Randrup J, *Phys. Rev. Lett.* **109** (2012) 212301
[9] Schaefer B J, Pawlowski J M and Wambach J, *Phys. Rev. D* **76** (2007) 074023
[10] Nahrgang M, Leupold S, Herold C and Bleicher M, *Phys. Rev. C* **84** (2011) 024912
[11] Herold C, Nahrgang M, Mishustin I and Bleicher M, *Phys. Rev. C* **87** (2013) 014907
[12] Nahrgang M, Herold C, Leupold S, Mishustin I and Bleicher M, *J. Phys. G* **40** (2013) 055108
[13] Herold C, Nahrgang M, Mishustin I and Bleicher M, arXiv:1304.5372 [nucl-th]