Second order chiral restoration phase transition at low temperatures in quarkyonic matter

L. Ya. Glozman and R. F. Wagenbrunn

Institute for Physics, Theoretical Physics branch, University of Graz, Universitätsplatz 5, A-8010 Graz, Austria

In this Addendum to our recent paper, Phys. Rev. D 77, 054027 (2008), we point out that a chiral restoration phase transition in a quarkyonic matter at low temperatures is of second order within a manifestly confining and chirally symmetric large \( N_c \) model. This result is qualitatively different as compared to NJL and NJL-like models that are not confining and might have some implications for the existence or nonexistence of the critical end point in the QCD phase diagram.

PACS numbers: 11.30.Rd, 12.38.Aw

In a recent paper \[1\] we studied a chiral restoration phase transition at finite density and zero temperature within the only known exactly solvable manifestly chirally \( SU(2)_L \times SU(2)_R \) symmetric and confining model in four dimensions \[2\]. It is assumed within this large \( N_c \) model that the only gluonic interaction is a linear confining potential of the Coulomb type. Then the chiral symmetry breaking can be obtained from the Schwinger-Dyson (gap) equation, while the color-singlet mesons are solutions of the Bethe-Salpeter equation. A single-quark Dirac operator is always infrared-divergent and hence a single quark is confined. All color-singlet observable quantities (quark condensate, hadron mass, ...) are infrared-finite and well defined. We have previously applied this model to study the chiral restoration in excited mesons \[3\].

A decisive feature of this model is that even above the chiral restoration point at a critical chemical potential the system is still in a confining mode and the only possible excitations are chirally invariant color-singlet hadrons (or meson-like color-singlet particle-hole excitations) \[1\] \[3\]. This is because in the large \( N_c \) limit there are no vacuum quark loops and no Debye screening of the confining gluon propagator at any chemical potential. The masses of the color-singlet and chirally-invariant hadrons increase with the chemical potential. This model represents a possible microscopic scenario for a recently proposed quarkyonic matter \[4\].

In this Addendum to our paper \[1\] we would like to point out that the chiral restoration phase transition at low temperatures is of second order. In order to see it explicitly we show in Fig. 1 the dependence of the chiral angle, \( \varphi_p \), which is a solution of a gap equation, on momentum \( p \) as well as on the quark Fermi-momentum \( p_f \) in a system with a finite chemical potential. The chiral restoration phase transition happens at the Fermi-momentum \( p_f^C = 0.109 \sqrt{\sigma} \), where \( \sigma \) is the string tension that supplies a dimensional scale in our task. At \( p_f = p_f^C \) the chiral angle takes its trivial value, \( \varphi_p = 0 \), and the chiral symmetry gets restored. An important feature is that the chiral angle approaches this trivial value continuously.

The chiral angle uniquely defines a dynamical mass \( M(p) \) of quarks, that is associated with the chiral symmetry breaking, as well as a quark condensate. Fig. 2 illustrates the continuous approach of the dynamical mass of quarks to its trivial value \( M(p) = 0 \) at the chiral restoration point. These two figures clearly demonstrate that the chiral restoration phase transition in the chiral limit
is of second order. The same can actually be seen from the continuous behavior of the quark condensate as well as the meson masses at the chiral restoration point (see Figs. 1 and 6 of Ref. [1]).

It is an interesting feature of this chirally symmetric and manifestly confining model. Indeed, it is well known that the Nambu and Jona-Lasinio models (NJL) exhibit a first order chiral restoration phase transition at low temperatures [2-4]. The same is true within the Polyakov-improved NJL models (PNJL) [5, 6], that simulate confinement in a statistical manner via a coupling of the non-confining NJL Hamiltonian with the Polyakov loop. It has been suggested very recently that this model also exhibits a chirally symmetric and confining quarkyonic phase [7, 8]. The chiral restoration phase transition is, however, of first order, like in the standard NJL model. This is in contrast to our manifestly confining model. One can speculate that this feature of the PNJL models is related to the absence of an explicit confinement of quarks.

Our results might have some implications for the existence/nonexistence of the critical end point, that is a subject of significant experimental and theoretical interest. The chiral restoration phase transition at zero density and large temperature is of second order in the chiral limit and becomes a crossover with realistic quark masses [10, 11]. Then, if at low temperatures and large density a chiral restoration phase transition is of first order, as suggested by the NJL-like and some other nonconfining models, there must exist a critical end point in the QCD phase diagram. If, however, a chiral restoration phase transition at low temperatures is of second order or a crossover, then there should be no critical end point (or, as an exotic possibility, there might be two critical end points).

Our results have been obtained within a large $N_c$ model, where the vacuum fermion loops as well as the Debye screening of the confining potential are absent. It would be interesting to see how persistent our results might be beyond the large $N_c$ limit.

Another caveat is that for the ground state of the system the implicit assumption is used that this ground state can be approximated as a quark Fermi gas, like in the NJL model. However, within the confining model the ground state is unlikely to be a pure Fermi gas of quarks. Going beyond a simple Fermi gas description of the ground state is a complicated task and should be considered as an important direction in the future.

Acknowledgements L.Ya.G. is thankful to T. Hatsuda, M. Huang, D. Kharzeev, L. McLerran, V. Miransky, R. Pisarski, B.-J. Schaefer, E. Shuryak, D. Son, J. Verbaarschot and J. Wambach for discussions and acknowledges support of the Austrian Science Fund through the grant P19168-N16.

[1] L. Ya. Glozman and R. F. Wagenbrunn, Phys. Rev. D 77, 054027 (2008).
[2] A. Le Yaouanc, L. Oliver, O. Pene, and J. C. Raynal, Phys. Rev. D 29, 1233 (1984); 31, 137 (1985); S. L. Adler and A. C. Davis, Nucl. Phys. B 244, 469 (1984); A. Kocic, Phys. Rev. D 33, 1785 (1986); R. Alkofer and P. A. Amundsen, Nucl. Phys. B 306, 305 (1988); P. Bicudo and J. E. Ribeiro, Phys. Rev. D 42 (1990) 1611; 42, 1625 (1990); P. J. A. Bicudo and A. V. Nefediev, Phys. Rev. D 68, 065021 (2003); F. J. Llanes-Estrada and S. R. Cotanch, Phys. Rev. Lett., 84, 1102 (2000); R. Alkofer, M. Kloker, A. Krassnigg, R. F. Wagenbrunn, Phys. Rev. Lett., 96, 022001 (2006).
[3] R. F. Wagenbrunn and L. Ya. Glozman, Phys. Lett. B 643, 98 (2006); R. F. Wagenbrunn and L. Ya. Glozman, Phys. Rev. D 75, 036007 (2007); L. Ya. Glozman, Phys. Rep. 444, 1 (2007).
[4] L. McLerran and R. D. Pisarski, Nucl. Phys. A 796, 83 (2007).
[5] L. Ya. Glozman, [arXiv:0803.1636] [hep-ph]
[6] M. Asakawa and K. Yazaki, Nucl. Phys. A 504, 668 (1089); B. Barducci et al, Phys. Lett. B 231, 463 (1989); Phys. Rev. D 41, 1610 (1990); S. P. Klevansky, Rev. Mod. Phys. 64, 646 (1992).
[7] P. N. Meisinger and M. C. Ogilvie, Phys. Lett. B 379, 163 (1996); K. Fukushima, Phys. Lett. B 591, 277 (2004); C. Ratti, M. A. Thaler, W. Weise, Phys. Rev. D 73, 014019 (2006); E. Megias, E. Ruiz Arriola, L.L. Saucedo, Phys. Rev. D 74, 114014 (2006); S. K. Ghosh et al, Phys. Rev. D 73, 114007 (2006); B. J. Schaefer, J. M. Pawlowski and J. Wambach, Phys. Rev. D 75, 074023 (2007); C. Sasaki, B. Friman and K. Redlich, Phys. Rev. D 75, 074013 (2007).
[8] K. Fukushima, [arXiv:0803.3318] [hep-ph]
[9] H. Abuki, R. Anglani, R. Gatto, G. Nardulli, M. Ruggeri, [arXiv:0805.1569] [hep-ph]
[10] F. Karsch, Plenary talk at Lattice 2007, PoS LATTICE2007:015,2006
[11] Z. Fodor et al, Plenary talk at Lattice 2007; LATTICE2007:189,2006