The QCD critical end point in the PNJL model

P. Costa$^{1,2(a)}$, C. A. de Sousa$^1$, M. C. Ruivo$^1$ and H. Hansen$^3$

$^1$ Departamento de Física, Universidade de Coimbra - P-3004-516 Coimbra, Portugal, EU
$^2$ E.S.T.G., Instituto Politécnico de Leiria - Morro do Lena-Alto do Vieiro, 2411-901 Leiria, Portugal, EU
$^3$ Université Lyon/UCBL, CNRS/IN2P3, IPNL - 4 rue E. Fermi, F-69622 Villeurbanne Cedex, France, EU

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Abstract – We investigate the role played by the Polyakov loop in the dynamics of the chiral phase transition in the framework of the so-called Polyakov-Nambu-Jona-Lasinio (PNJL) model in the $SU(2)$ sector. We present the phase diagram where the inclusion of the Polyakov loop moves the critical points to higher temperatures, compared with the Nambu-Jona-Lasinio model results. The critical properties of physical observables, such as the baryon number susceptibility and the specific heat, are analyzed in the vicinity of the critical end point, with special focus on their critical exponents. The results with the PNJL model are closer to lattice results and we also recover the universal behavior of the critical exponents of both the baryon susceptibility and the specific heat.

Confinement and chiral symmetry breaking are two of the most important features of quantum chromodynamics (QCD). Chiral models like the Nambu-Jona-Lasinio (NJL) model have been successful in explaining the dynamics of spontaneous breaking of chiral symmetry and its restoration at high temperatures and densities/chemical potentials. Recently, this and other types of models, together with an intense experimental activity, are underway to construct the phase diagram of QCD.

Results with two massless quarks in QCD show that, at high temperature, the phase transition associated to restoration of chiral symmetry is second order and belongs to the universality class of O(4) spin models in three dimensions [1]. With small quark masses, the second-order phase transition is replaced by a smooth crossover, a picture which is consistent with lattice simulations [2]. Various results from QCD-inspired models indicate (e.g., refs. [3,4]) that at low temperatures the transition may be first order for large values of the chemical potential. This suggests that the first-order transition line may end when the temperature increases, the phase diagram thus exhibiting a critical end point (CEP) [5–7] that can be detected via enhanced critical fluctuations in heavy-ion reactions [8,9]. At the CEP, the transition is second order and belongs to the Ising universality class [10]. In the chiral limit a tricritical point (TCP) is found in the phase diagram, separating the second-order transition line from the first-order one.

Recent developments in lattice QCD [11] indicate that the CEP is likely to be localized by a new generation of experiments with relativistic nuclei (CBM experiment at FAIR), suggesting to explore the range of baryon number chemical potential $\mu_B = 100$–$500$ MeV.

In a previous work [7], in the framework of the Nambu-Jona-Lasinio (NJL) model, we studied the phase diagram, focusing our attention on the CEP and the physics near it, through the behavior of the baryon number susceptibility and the specific heat.

In this work we study thermodynamic properties of strongly interacting matter using the Polyakov-Nambu-Jona-Lasinio (PNJL) model. This extended model, first implemented in ref. [12], provides a simple framework that couples the chiral and the confinement order parameters. The NJL model describes interactions between constituent quarks, giving the correct chiral properties; static gluonic degrees of freedom are then introduced in the NJL Lagrangian through an effective gluon potential in terms of Polyakov loops [12–14] with the aim of taking into account features of both chiral symmetry breaking and deconfinement. The coupling of the quarks to the Polyakov loop leads to the reduction of the weight of the quarks degrees of freedom as the critical temperature is

$^{(a)}$E-mail: pcosta@teor.fis.uc.pt
approached from above, which is interpreted as a manifestation of confinement and is essential to reproduce lattice results. We emphasize that the reduction of the weight of the quark degrees of freedom might also have an important role for the critical behavior.

This effect should be more visible in the temperature domain, which can be explained by the attractive interactions between the quarks and the effective gluon field which shifts the chiral phase transition temperature to high values, allowing for a stronger first-order phase transition.

Hence it is demanding to use this improved NJL model to investigate relevant thermodynamical quantities such as the CEP and the TCP.

Our main goal is to locate the critical end point in the PNJL model [15] and confront the results with the NJL one and universality arguments. Based on the fact that the CEP is a genuine thermodynamic singularity, being considered a second-order critical point, response functions like the specific heat and susceptibilities can provide relevant signatures for phase transitions. We notice that susceptibilities in general are related to fluctuations through the fluctuation dissipation theorem, allowing to observe signals of phase transitions in heavy-ion reactions [16,17].

The Lagrangian of the SU(2) ⊗ SU(2) quark model with explicit chiral symmetry breaking where the quarks couple to a (spatially constant) temporal background gauge field (represented in term of Polyakov loops) is given by [14,18]

\[
\mathcal{L}_{PNJL} = \bar{q}(i\gamma^\mu D_\mu - m)q + \frac{1}{2} g S \left( (\bar{q}q)^2 + (\bar{q}i\gamma_5 T q)^2 \right) - \mathcal{U}(\Phi[A], \bar{\Phi}[A]; T),
\]

(1)

The quark fields \( q = (u, d) \) are defined in Dirac and color fields, respectively, with two flavors, \( N_f = 2 \) and three colors, \( N_c = 3 \), and \( m = \text{diag}(m^0_u, m^0_d) \) is the current quark mass matrix.

The quarks are coupled to the gauge sector via the covariant derivative \( D^\mu = \partial^\mu - iA^\mu \). The strong coupling constant \( g_{\text{strong}} \) has been absorbed in the definition of \( A^\mu = g_{\text{strong}}A_\mu^a(x)\tau^a \), where \( A_\mu^a \) is the SU(3) gauge field and \( \lambda^a \) are the Gell-Mann matrices. Beside, in the Polyakov gauge and at finite temperature \( A^\mu = \delta^\mu_0 A^0 = -i\delta^\mu_0 A^4 \).

The Polyakov loop \( \Phi \) (the order parameter of \( Z_3 \) symmetric/broken phase transition in pure gauge) is the trace of the Polyakov line defined by \( \Phi = \frac{1}{N_f} \langle \langle \mathcal{P} \exp i \int_0^\beta d\tau A_4(x, \tau) \rangle \rangle \).

The pure gauge sector is described by an effective potential \( \mathcal{U}(\Phi[A], \bar{\Phi}[A]; T) \) chosen to reproduce at the mean-field level the results obtained in lattice calculations:

\[
\frac{\mathcal{U}(\Phi, \bar{\Phi}; T)}{T^4} = -\frac{b_2(T)}{2} (\Phi + \bar{\Phi}) - \frac{b_3}{6} (\Phi^3 + \bar{\Phi}^3) - \frac{b_4}{4} (\Phi \bar{\Phi})^2 ,
\]

(2)

where

\[
b_2(T) = a_0 + a_1 \left( \frac{T_0}{T} \right) + a_2 \left( \frac{T_0}{T} \right)^2 + a_3 \left( \frac{T_0}{T} \right)^3
\]

and \( a_0 = 6.75, a_1 = -1.95, a_2 = 2.625, a_3 = -7.44, b_3 = 0.75, b_4 = 7.5 \), and \( T_0 = 270 \text{ MeV} \).

The parameters of the pure NJL sector are fixed at zero temperature as in [6]: the three-momentum cutoff used to regularize all the integrals is \( \Lambda = 590 \text{ MeV} \), \( m^0_u = m^0_d = 6 \text{ MeV} \), and \( g_S A^2 = 2.435 \). They yield \( M_{vac} = 400 \text{ MeV} \), \( m_o = 140.2 \text{ MeV} \), \( f_r = 92.6 \text{ MeV} \), and \( \langle \bar{u}u \rangle^{1/3} = (-241.5 \text{ MeV})^3 \).

Finally with \( E_2^2 = p^2 + M^2 \), the \( SU(N_f = 2) \) PNJL grand potential is given by [14,19]

\[
\text{see eq. (3) above}
\]

We notice explicitly that at \( T = 0 \) the Polyakov loop and the quark sector decouple.

The baryon number susceptibility and the specific heat are the response of the baryon number density \( \rho_B(T, \mu) \) and the entropy \( S(T, \mu) \) to an infinitesimal variation of the quark chemical potential \( \mu \) and temperature given, respectively, by

\[
\chi_q = \left( \frac{\partial \bar{\rho}_q}{\partial \mu} \right)_T \quad \text{and} \quad C = \frac{T}{V} \left( \frac{\partial S}{\partial T} \right)_{\mu}.
\]

(4)

The baryon number density is given by \( \rho_q = \sum_q n_q(\mu, T) \), where \( n_q(\mu, T) \) and \( \bar{n}_q(\mu, T) \) are the occupation numbers modified by the Polyakov loop [19].

The PNJL thermodynamic potential is an effective potential depending on three parameters: \( M, \Phi \), and \( \Phi \). These parameters are not independent (nor the corresponding phase transitions) since they should verify the
The QCD critical endpoint in the PNJL model

Fig. 1: (Color online) Left panel: The phase diagram in the PNJL and NJL models. Right panel: The size of the critical region is plotted for $\chi_q/\chi_{q,\text{free}} = 2, 3, 5$.

The size of the critical region around the CEP can be found by calculating the baryon number susceptibility, the specific heat, and their critical behaviors. The size of this critical region is important for future searches of the CEP in heavy-ion collisions [17]. To estimate the critical region around the CEP, we calculate the dimensionless ratio $\chi_q/\chi_{q,\text{free}}$, where $\chi_q$ is obtained taking the chiral limit $\mu = 0$. The right panel of fig. 1 shows a contour plot for three fixed ratios ($\chi_q/\chi_{q,\text{free}} = 2.0, 3.0, 5.0$) in the phase diagram around the CEP where we notice an elongation of the region where $\chi_q$ is enhanced, in the direction parallel to the first-order transition line. We also observe that the critical region is heavily stretched in the direction of the crossover transition line as shown in fig. 1.

The elongation of the critical region in the $(T, \mu)$-plane, along the critical line, is larger in the PNJL model (see $\chi_q/\chi_{q,\text{free}} = 2.0$ in the right panel of fig. 1). It means that the divergence of the correlation length at the CEP affects the phase diagram quite far from the CEP and that a careful analysis including effects beyond the mean field needs to be done [20].

As seen in [19] (fig. 7), one of the main effects of the Polyakov loop is to shorten the temperature range where the crossover occurs (at $\mu = 0$, the crossover occurs within a range of 150 MeV for the NJL model and within 115 MeV for the PNJL one), thus resulting in higher baryonic susceptibilities even far from the CEP. This effect is driven by the fact that the one- and two-quark Boltzmann factors are controlled by a factor proportional to $\Phi$: at small temperature, $\Phi \simeq 0$, results in a suppression of these contributions. The thermal bath being then only produced via the three-quark Boltzmann factor, our physical interpretation is that the bath is colorless, quarks being produced only in triplet necessarily colorless in the average because of $\Phi$ being the order parameter of $Z_3$ in this effective theory, and $\Phi \simeq 0$ indicates a partial restoration of the color symmetry. When the temperature increases, $\Phi$ goes quickly to 1, resulting in a (partial) restoration of

mean-field equations $\partial \Omega / \partial M = 0$ and $\partial \Omega / \partial \Phi = \partial \Omega / \partial \bar{\Phi} = 0$. With the two last equations, one can compute $\Phi$ and $\bar{\Phi}$ as functions of $M$ for any value of $T$ and $\mu$. Hence, we consider that the thermodynamic potential is an effective potential depending only on $M$.

A common feature shared by NJL and PNJL models is that the thermodynamic potential may have two degenerate minima at which two phases have equal pressure and chemical potential and can coexist according to the Gibbs criterion. In fact, this pattern is characteristic of a first-order phase transition: the two minima correspond, respectively, to the phases of broken and restored symmetry. The quark condensate can be identified with the order parameter whose values allow to distinguish the two coexisting phases. As the temperature increases, the first-order transition persists up to the CEP. At the CEP, the chiral transition becomes of the second order. For temperatures above the CEP, the thermodynamic potential has only one minimum and the transition is washed out: a smooth crossover takes place.

So, while from a qualitative point of view the results of both models are similar, they differ quantitatively in two aspects: in the PNJL, the CEP and TCP are pushed to higher values of the temperature and the size of the critical region is larger. Let us now analyze those results in more detail.

In the left panel of fig. 1, we plot the phase diagram for both PNJL and NJL models. In the PNJL (NJL) model, the CEP is localized at $T_{\text{CEP}}^{\text{PNJL}} = 169.11$ (79.92) MeV and $\mu_{\text{CEP}}^{\text{PNJL}} = 321.32$ (331.72) MeV ($\rho_{q,\text{CEP}}^{\text{PNJL}} = 2.76$ (2.95)$\rho_0$). In the PNJL (NJL) model, the TCP is located at $T_{\text{TCP}}^{\text{PNJL}} = 207.66$ (112.08) MeV and $\mu_{\text{TCP}}^{\text{PNJL}} = 270.80$ (286.05) MeV. We remark that our values for the CEP in the PNJL model are closer to the lattice results of [11] and the main change, with regard to NJL values, is in $T_{\text{CEP}}^{\text{PNJL}}$ and $T_{\text{TCP}}^{\text{PNJL}}$, a result that seems natural since the effects of the inclusion of the Polyakov loop are expected to be more relevant in the domain of the temperature.
the chiral symmetry that occurs in a shorter temperature range. In fact, the most striking difference between NJL and PNJL models is a faster variation with temperature, around any characteristic critical temperature, of the PNJL results.

The crossover taking place in a smaller temperature range can be interpreted as a crossover transition closer to a second-order one than in the NJL model. This “faster” crossover may also explain the elongation of the critical region compared to the NJL one giving rise to a greater correlation length even far from the CEP.

With this indication of the important role of the entanglement of the chiral and the Polyakov loop dynamics on the critical behavior of the QCD phase diagram, it is mandatory to investigate the behavior of $\chi_q$ and $C$ in the vicinity of the CEP and their critical exponents, in the framework of the PNJL model. For comparison purposes with the NJL model and the universality/mean-field predictions, the calculated critical exponents at CEP and the TCP are presented in table 1, and will be discussed in the sequel.

The phenomenological relevance of fluctuations in the finite temperature and chemical potential around the CEP/TCP of QCD has been recognized by several authors. If the critical region of the CEP is small, it is expected that most of the fluctuations associated with the CEP will come from the mean-field region around the CEP [9].

In the left panel of fig. 2, $\chi_q$ is plotted as a function of $\mu$ for three different temperatures around the CEP. For temperatures below $T^{CEP}$, we have a first-order phase transition and, consequently, $\chi_q$ has a discontinuity. For $T = T^{CEP}$, the slope of the baryon number density tends to infinity at $\mu = \mu^{CEP}$, which implies a diverging $\chi_q$. For temperatures above $T^{CEP}$, in the crossover region, the discontinuity of $\chi_q$ disappears at the transition line.

A similar behavior is found for the specific heat as a function of temperature for three different chemical potentials around the CEP, as we can see from the right panel of fig. 2.

These behaviors of $\chi_q$ and $C$ are qualitatively similar to those obtained in the $SU(2)$ NJL model [7]. As we have already seen, the baryon number susceptibility, $\chi_q$, and the specific heat, $C$, diverge at $T = T^{CEP}$ and $\mu = \mu^{CEP}$, respectively [7,9]. In order to make this statement more precise, we will focus on the values of the critical exponents, in our case $\epsilon$ and $\alpha$ are the critical exponents of $\chi_q$ and $C$, respectively. These critical exponents will be determined by finding two directions, temperature-like and magnetic-field-like, in the $(T, \mu)$-plane near the CEP, because, as pointed out in [21], the form of the divergence depends on the route which is chosen to approach the critical end point.

To study the critical exponents for the baryon number susceptibility (eq. (4)), we will start with a path parallel to the $\mu$-axis in the $(T, \mu)$-plane, from lower $\mu$ toward the critical $\mu^{CEP} = 321.32$ MeV, at fixed temperature $T^{CEP} = 169.11$ MeV. In fig. 3, we plot $\chi_q$ as a function of $\mu$ close to the CEP. Using a linear logarithmic fit

$$\ln \chi_q = -\epsilon \ln |\mu - \mu^{CEP}| + c_1,$$

where the term $c_1$ is independent of $\mu$, we obtain $\epsilon = 0.66 \pm 0.01$, which is consistent with the mean-field theory prediction $\epsilon = 2/3$.

We also study the baryon number susceptibility from higher $\mu$ toward the critical $\mu^{CEP}$. The logarithmic fit used now is $\ln \chi_q = -\epsilon' \ln |\mu - \mu^{CEP}| + c_1'$. Our result shows that $\epsilon' = 0.69 \pm 0.02 \approx \epsilon$. This means that the size of the region we observe is approximately the same independently of the direction we choose for the path parallel to the $\mu$-axis. These critical exponents are presented in table 1, where we can see that the critical exponents for the baryon number susceptibility are approximately the same for both, PNJL and NJL models, and are consistent with the mean-field theory prediction $\epsilon = 2/3$. 

![Fig. 2: (Color online) Left panel: baryon number susceptibility as a function of $\mu$ for different temperatures around the CEP in PNJL model: $T^{CEP} = 169.11$ MeV and $T = T^{CEP} \pm 10$ MeV. Right panel: specific heat as a function of $T$ for different values of $\mu$ around the CEP: $\mu^{CEP} = 321.32$ MeV and $\mu = \mu^{CEP} \pm 10$ MeV.](image_url)
The QCD critical end point in the PNJL model

Table 1: (Color online) Critical exponents (CE): the arrow → (●) indicates the path in the $\mu$ ($T$)-direction to the CEP (TCP) for $\mu < \mu^{CEP} (T < T^{CEP})$.

| Quantity | CE/path | PNJL | NJL | Universality |
|----------|---------|------|-----|--------------|
| $\chi_q$ | $\epsilon$ | 0.66 ± 0.01 | 0.66 ± 0.01 | 2/3 |
| $\epsilon'$ | $\gamma_q$ | 0.69 ± 0.02 | 0.66 ± 0.01 | 2/3 |
| $\alpha$ | $\alpha'/\epsilon$ | 0.51 ± 0.01 | 0.51 ± 0.01 | 1/2 |

On the other hand, in the chiral limit (where the CEP becomes a TCP), it is found that the critical exponent for $\chi_q$ has the value $\gamma_q = 0.51 \pm 0.01$, for both, PNJL and NJL models. Again, these results are in agreement with the mean-field value ($\gamma_q = 1/2$).

Now, paying attention to the specific heat around the CEP, we have used a path parallel to the $T$-axis in the $(T,\mu)$-plane from lower (higher) $T$ toward the critical $T^{CEP} = 169.11$ MeV at fixed $\mu^{CEP} = 321.32$ MeV. In fig. 3 (right panel), we plot $C$ as a function of $T$ close to the CEP in a logarithmic scale. We see that for the region $T < T^{CEP}$, we have $\alpha = 0.63 \pm 0.02$\(^1\). Contrarily to what happens in the NJL model (see table 1 and ref. [7]), this value of $\alpha$ is closer to the one suggested by universality arguments in [9].

We also observe, as in ref. [7], that in PNJL (NJL) model, for the region $T < T^{CEP}$ we have a slope of data points that changes for values of $|T - T^{CEP}|$ around 0.3 MeV. We have fitted the data for $|T - T^{CEP}| < 0.3$ MeV and $|T - T^{CEP}| > 0.3$ MeV separately and obtained, respectively, the critical exponent $\alpha = 0.63 \pm 0.02$ ($\alpha = 0.59 \pm 0.01$) and $\alpha_1 = 0.53 \pm 0.01$ ($\alpha_{1\uparrow} = 0.45 \pm 0.01$), which have a linear behavior for several orders of magnitude (see table 1). As pointed out in [9], this change of the exponent can be interpreted as a crossover of different universality classes, with the CEP being affected by the TCP.

In both models, the influence of the TCP is stronger in the specific heat rather than in the baryon number susceptibility: the closest distances between the TCP and the CEP in both phase diagrams occur in the $T$-direction ($(T^{TCP} - T^{CEP}) < (\mu^{CEP} - \mu^{TCP})$). When the CEP is approached from above the trivial exponent $\alpha' = 0.69$ (for both models) is obtained.

Let us now analyze the behavior of the specific heat around the TCP. As shown in table 1, we find a nontrivial critical exponent $\alpha = 0.40 \pm 0.02$ only for the NJL model while for the PNJL model $\alpha = 0.50 \pm 0.01$.

In this work, we have considered an extension of the NJL model which couples chiral and confinement-like order parameters. We have found that our model in general reproduces important features of the QCD phase diagram as the location of the CEP/TCP. In addition, these results confirm the general idea that, in contrast to the NJL model, the PNJL model provides a quantitative description of QCD thermodynamics near critical points. In the PNJL model, the crossover taking place in a smaller $T$ range can be interpreted as a crossover transition closer to a second-order one than in the NJL model. This “faster” crossover may explain the elongation of the critical region compared to the NJL one giving rise to a greater correlation length even far from the CEP. We

\(^1\)We use the linear logarithmic fit in $C = -\alpha \ln|T - T^{CEP}| + c_2$, where the term $c_2$ is independent of $T$. 

Fig. 3: (Color online) Left panel: baryon number susceptibility as a function of $|\mu - \mu^{CEP}|$ at the fixed $T = T^{CEP}$. Right panel: specific heat as a function of $|T - T^{CEP}|$ at the fixed $\mu = \mu^{CEP}$.
have also studied the baryon number susceptibility and the specific heat around the CEP which are related with event-by-event fluctuations of $\mu$ or $T$ in heavy-ion collisions. An important observation is that, in the PNJL model, the obtained critical exponents are consistent with the mean-field values, both for the baryon number susceptibility and the specific heat, while for the NJL this is only true for the baryonic susceptibility, since for the specific heat $\alpha$ is different from $\epsilon$.

As the CEP lies in the region expected to be probed by heavy-ion experiments, it would be interesting to find an experimental signature of such a point. Near the critical point, and in particular in the path we choose to study the critical exponents of the specific heat, there is a possibility of the spinodal decomposition in the first-order phase transition. Hence, the competition between features of the first- and second-order phase transitions in the mixed phase can allow for nontrivial effects, such as the above referred ones and to which there is no information from heavy-ion collisions. Our numerical results that also include the chemical potential can be relevant to this purpose.

In conclusion, the results with the PNJL model are closer to lattice results and we also recover the universal behavior of the critical exponents of both the baryon susceptibility and the specific heat. The PNJL model discussed here, allowing for finite dynamical quark masses, can provide a convenient tool to study the QCD phase diagram; it allows to establish a convenient link between the lattice results, and the NJL model itself where gluonic degrees of freedom are missing.

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REFERENCES

[1] Rajagopal K., Nucl. Phys. A, 661 (1999) 150.
[2] Laermann E., Nucl. Phys. Proc. Suppl., 63 (1998) 114.
[3] Stephanov M. A., Phys. Rev. Lett., 76 (1996) 4472.
[4] Barducci A. et al., Phys. Lett. B, 231 (1989) 463; Phys. Rev. D, 41 (1990) 1610.
[5] Schwarz T. M., Klevansky S. P. and Papp G., Phys. Rev. C, 60 (1999) 055205.
[6] Frank M., Buballa M. and Oertel M., Phys. Lett. B, 562 (2003) 221.
[7] Costa P., de Sousa C. A., Ruiz Arriola E., Phys. Lett. B, 647 (2007) 431; Costa P., Ruivo M. C. and de Sousa C. A., Phys. Rev. D, 77 (2008) 096001.
[8] Stephanov M., Rajagopal K. and Shuryak E., Phys. Rev. Lett., 81 (1998) 4816.
[9] Hatta Y. and Ikeda T., Phys. Rev. D, 67 (2003) 014028; Schaefer B.-J. and Wambach J., Phys. Rev. D, 75 (2007) 085015.
[10] Berges J. and Rajagopal K., Nucl. Phys. B, 538 (1999) 215; Halasz M. A. et al., Phys. Rev. D, 58 (1998) 096007.
[11] Karsch F., AIP Conf. Proc., 842 (2006) 20; Fodor Z. and Katz S. D., J. High Energy Phys., 0204 (2004) 050.
[12] Meisinger P. N. and Ogilve M. C., Phys. Lett. B, 379 (1996) 163.
[13] Megias E., Ruiz Arriola E. and Salcedo L. L., Phys. Rev. D, 74 (2006) 065005; 114014; Ciminale M. et al., Phys. Rev. D, 77 (2008) 054023.
[14] Ratti C., Thaler M. A. and Weise W., Phys. Rev. D, 73 (2006) 014019; Fukushima K., Phys. Lett. B, 591 (2004) 277.
[15] Kashiwa K., Kouno H., Matsuzaki M. and Yahiro M., Phys. Lett. B, 662 (2008) 26; Schaefer B.-J., Pawlowski J. M. and Wambach J., Phys. Rev. D, 76 (2007) 074023; Sasaki C., Friman B. and Redlich K., Phys. Rev. D, 75 (2007) 074013; Rößner S., Ratti C. and Weise W., Phys. Rev. D, 75 (2007) 034007.
[16] Appelshauser H. and Sako H. the CERES Collaboration, Nucl. Phys. A, 752 (2005) 394; PHENIX Collaboration (Mitchell J. T.), J. Phys. Conf. Ser., 27 (2005) 88.
[17] Nonaka C. and Asakawa M., Phys. Rev. C, 71 (2005) 044904.
[18] Pisarski R. D., Phys. Rev. D, 62 (2000) 111501R; hep-ph/0203271.
[19] Hansen H., Alberico W., Beraudo A., Molinari A., Nardi M. and Ratti C., Phys. Rev. D, 75 (2007) 065004.
[20] Rößner S., Hell T., Ratti C. and Weise W., Nucl. Phys. A, 814 (2008) 118.
[21] Griffiths R. B. and Wheeler J., Phys. Rev. A, 2 (1970) 1047.

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