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Dilepton production from ρ-mesons in a quark-gluon plasma

M.H. Thoma[1,2,a], S. Leupold[1], U. Mosel[1]

[1] Institut für Theoretische Physik, Universität Giessen, 35392 Giessen, Germany
[2] ECT*, Villa Tambosi, Strada delle Tabarelle 286, 38050 Villazzano (Trento), Italy

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Abstract. Assuming that ρ-mesons exist in a quark-gluon plasma at temperatures close to the QCD phase transition, we calculate the dilepton production rate from q-$\bar{q}$ annihilation via a ρ -meson state using Vector Meson Dominance. The result is compared to the rates from direct q-$\bar{q}$ annihilation and from π^+-π^- annihilation. Furthermore we discuss the suppression of low mass dileptons if the quarks assume an effective mass in the quark-gluon plasma.

PACS. 25.75.-q Relativistic heavy-ion collisions – 12.38.Mh Quark-gluon plasma – 14.40.Cs Other mesons with S = C = 0, mass <2.5 GeV

Dileptons are one of the most promising signatures for medium effects in the fireball in relativistic heavy ion collisions. At SPS there is an indication for an enhancement of the dilepton production at invariant masses between 200 and 800 MeV compared to all known sources neglecting medium effects [1]. Possible explanations of this enhancement are the broadening of the width of the ρ -meson by scattering in the medium and the reduction of the ρ-mass by the on-set of the restoration of chiral symmetry [2].

Furthermore dileptons may indicate the formation of a quark-gluon plasma (QGP) phase, in particular at RHIC and LHC [3]. In a perturbative calculation the lowest order contribution to the dilepton production from the QGP comes from the direct q-$\bar{q}$ annihilation into a virtual photon (Born term) [4]. At high invariant masses M this contribution is believed to dominate. For low invariant masses of the order of gT, on the other hand, α_s-corrections become important [5–7]. Close to the critical temperature g might be as large as 6 [8]. Hence α_s-corrections could be important even at invariant masses up to $M = 1$ GeV or larger.

In order to avoid infrared singularities and gauge dependent results, HTL resummed propagators and vertices [9] have to be used in perturbative calculations of the dilepton production from the QGP [6,7]. The dilepton rate follows from the imaginary part of the finite temperature photon self energy. Unfortunately, it turned out that 2- and probably even 3-loop contributions [10] to the photon self energy are of the same order in α_s as the 1-loop contribution and exceed the latter one even by factors of 3 and more. Thus the application of finite temperature perturbation theory suffers from the facts that higher order loop corrections are important even in the weak coupling limit and moreover that realistic values of the strong coupling constant are not small.

The importance of non-perturbative effects in a QGP at temperatures within reach of heavy ion collisions has been realized in QCD lattice calculations, where for example the equation of state, condensates, and hadronic correlators below and above the phase transition have been studied [11].

It has been noticed that the equation of state found on the lattice can be described perfectly by an ideal gas of massive quarks and gluons (quasiparticles), where an effective, temperature dependent mass of the order gT for the partons has been introduced [8]. The effective quark and gluon mass at temperatures around the critical one T_c is of the order 0.5 GeV. This will lead to a complete suppression of dileptons from the QGP below about $M = 1$ GeV as we will discuss at the end of the present paper.

The existence of a gluon condensate above T_c can be used to construct an effective quark propagator in the QGP [12]. The quark dispersion relation following from this propagator has two branches, of which one shows a minimum at finite momentum. This leads to an interesting structure of the dilepton spectrum, exhibiting peaks (Van Hove singularities) and gaps [6,13].

In the present paper we want to study the consequences of the existence of hadronic states in the QGP at temperatures around T_c as indicated by lattice calculations [14] and the gauged linear sigma model [15] on the dilepton production from the QGP. This temperature regime plays an important role in relativistic heavy ion collisons as the mixed phase with $T = T_c$ may exist for a long period in the fireball and might contribute significantly to

[a] Heisenberg fellow

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the photon and dilepton production (see e.g. [16]). In particular, we consider a possible co-existence of quarks and ρ-mesonic states in the QGP, because ρ-mesons are an important source for dileptons. Although one expects that ρ-mesons vanish above T_c quickly with increasing temperature [14], they might be present at temperatures around T_c [15].

We assume a simple coupling of the quarks to ρ-mesonic states in the QGP. Besides by direct q-$\bar{q}$ annihilation electron-positron pairs can then be produced by the process $q\bar{q} \to \rho \to e^+e^-$, where the decay of the ρ will be described using Vector Meson Dominance (VMD). The aim of this investigation is to compare our results for this particular mechanism to the rates following from direct annihilation of quarks (Born term) and from $\pi^+\pi^-$ annihilation using VMD. Furthermore we assume a vanishing baryon density (zero quark chemical potential), equilibrium distributions for the quarks and ρ-mesons, and the bare mass $m_\rho = 770$ MeV for the ρ-meson.

For the interaction of quarks with ρ-mesons we adopt the following Lagrangian:

$$\mathcal{L} = -\frac{1}{4}\,\rho^a_{\mu\nu}\,\rho^{\mu\nu}_a + \frac{1}{2}\,m^2_\rho\,\rho^a_\mu\,\rho^\mu_a + \bar{q}\,\left(i\,\gamma_\mu\,\partial^\mu - m_q + g\,\gamma^\mu\,\frac{\tau_a}{2}\,\rho^a_\mu\right) q\,, \quad (1)$$

where $\rho^a_{\mu\nu} = \partial_\mu\rho^a_\nu - \partial_\nu\rho^a_\mu$ with the ρ-meson field ρ^a_μ and the quark field q. Here a denotes the isospin or flavor index and τ_a the corresponding isospin matrix. For the quark mass m_q we will investigate two cases, namely a vanishing bare mass and an effective quark mass $m_q = 0.5$ GeV.

To get an idea about the size of the coupling constant g describing the strength of the quark-ρ coupling we proceed in the following way: By integrating out the ρ-mesons from Lagrangian (1) one obtains in lowest order of the derivative expansion the four-quark interaction term

$$\mathcal{L}_{4-\mathrm{quark}} = -\frac{1}{2}\frac{g^2}{m^2_\rho}\left(\bar{q}\gamma_\mu\frac{\tau_a}{2}q\right)^2, \quad (2)$$

which might be compared with the corresponding interaction term from the extended NJL model $-G_2\,(\bar{q}\gamma_\mu\tau_a q)^2$ [17,18]. This suggests the identification $g = \sqrt{8m^2_\rho G_2}$ which serves to determine g by taking G_2 from the literature [17,18]: $g \approx 5-6$. In the following we choose $g = 6$. The same result for g can be obtained by simply assuming that the ρ-meson couples in a universal way to nucleons, pions and quarks. From [2] (fourth reference) as well as from [19] one would get $g \approx 6.1$. Of course, it is not clear whether our choice for mass and coupling of a vector meson excitation in a hot QGP is justified. However, in lack of any first principle calculation we decided to choose values as suggested by measurements or calculations for the vacuum case.

The dilepton production rate (here for massless e^+-e^- pairs) can be calculated from the imaginary part of the photon self energy according to [19]

$$\frac{dN}{d^4x d^4p} = -\frac{\alpha}{12\pi^4}\,\frac{1}{e^{E/T}-1}\,\frac{Im\Pi^\mu_\mu(P)}{M^2}, \quad (3)$$

where $\alpha = e^2/4\pi$. Here we use the notation $P = (E, \boldsymbol{p})$ and $p = |\boldsymbol{p}|$.

Using VMD the photon self energy is related to the ρ^0-meson propagator by

$$Im\Pi^\mu_\mu(P) = \frac{e^2}{g^2_\rho}\,m^4_\rho\,ImD^\mu_\mu(P), \quad (4)$$

where $g_\rho = 6.07$ [19]. The most general ansatz for the trace of the imaginary part of the full ρ-meson propagator reads

$$ImD^\mu_\mu(P) = -\frac{ImF}{(M^2 - m^2_\rho - ReF)^2 + (ImF)^2} - \frac{2\,ImG}{(M^2 - m^2_\rho - ReG)^2 + (ImG)^2}, \quad (5)$$

where $F(p_0, p)$ and $G(p_0, p)$ are the longitudinal and transverse parts of the ρ^0-self energy. To lowest order the ρ^0-self energy is calculated from the one-loop diagram containing a quark loop. It is given by the one-loop photon self energy containing an electron loop multiplied by an factor 3/2, where the factor 3 comes from the number of quark colors in the loop and the factor 1/2 from the flavor coefficient $\mathrm{tr}(\tau^2_0/4)$, which counts the number of quark flavors (u and d). Furthermore the electron charge e has to be replaced by g.

We restrict ourselves to the calculation of the dilepton rate from the one-loop photon self energy taking into account only bare quark propagators as we want to compare our results with the Born rate, where also only bare quarks have been considered in a one-loop approximation. Further effects, as the ones considered in [6] using HTL resummed quark propagators and vertices, which are important at small invariant masses, are beyond the scope of the present work. It should be noted also that in our model HTL diagrams are suppressed by the large mass of the ρ-meson. Therefore these effects are of higher order in g even at small invariant mass in contrast to QCD calculations.

The one-loop photon self energy at finite temperature can be calculated analytically in the high temperature or equivalently hard thermal loop limit [9,20]. However, in this approximation there is no imaginary part for time-like photons, $M^2 = p_0^2 - p^2 > 0$, resulting in a vanishing dilepton production. Going beyond the hard thermal loop approximation, integral expressions for the matter part of the one-loop photon self energy can be derived [21]:

$$ReF = \frac{3}{2\pi^2}\,g^2\,\frac{M^2}{p^2}\int_0^\infty dk\,k\,n_F(\omega_k)\,\Bigg[-2\frac{k}{\omega_k} + \frac{4\omega^2_k + M^2}{4p\omega_k}\,\ln\left|\frac{(2pk+M^2)^2 - 4p^2_0\omega^2_k}{(2pk-M^2)^2 - 4p^2_0\omega^2_k}\right| + \frac{p_0}{p}\,\ln\left|\frac{M^4 - 4(p_0\omega_k + pk)^2}{M^4 - 4(p_0\omega_k - pk)^2}\right|\Bigg],$$

$$
\begin{aligned}
ImF &= \frac{3}{2\pi}\, g^2\, \frac{M^2}{p^3}\, \left[\int_{k_-}^{k_+} dk\, k\, n_F(\omega_k) \left(p_0 - \omega_k - \frac{M^2}{4\omega_k} \right) \right. \\
&\quad \left. -2\, p_0\, \theta(-M^2) \int_{k_-}^{\infty} dk\, k\, n_F(\omega_k) \right], \\
ReG &= \frac{3}{2\pi^2}\, g^2 \int_0^{\infty} dk\, \frac{k^2}{\omega_k}\, n_F(\omega_k) \\
&\quad \times \left[2 + \frac{M^2}{p^2} - \left(\frac{\omega_k^2 M^2}{2p^3 k} + \frac{M^2}{4pk} + \frac{M^4}{8p^3 k} + \frac{m_q^2}{2pk} \right) \right. \\
&\quad \times \ln \left| \frac{(2pk + M^2)^2 - 4p_0^2 \omega_k^2}{(2pk - M^2)^2 - 4p_0^2 \omega_k^2} \right| - \frac{p_0 M^2 \omega_k}{2p^3 k} \\
&\quad \left. \times \ln \left| \frac{M^4 - 4(p_0 \omega_k + pk)^2}{M^4 - 4(p_0 \omega_k - pk)^2} \right| \right], \\
ImG &= \frac{3}{4\pi}\, g^2\, \frac{1}{p} \left[\int_{k_-}^{k_+} dk\, k\, n_F(k) \right. \\
&\quad \times \left(-k + \frac{p_0^2}{p^2} \omega_k + \frac{M^2}{2\omega_k} + \frac{M^4}{4p^2 \omega_k} - \frac{m_q^2}{\omega_k} - \frac{p_0 M^2}{p^2} \right) \\
&\quad \left. + \frac{p_0 M^2}{p^2}\, \theta(-M^2) \int_{k_-}^{\infty} dk\, k\, n_F(\omega_k) \right], \qquad (6)
\end{aligned}
$$

where $\omega_k^2 = k^2 + m_q^2$, $n_F(\omega_k) = 1/[\exp(\omega_k/T) + 1]$, and

$$
\begin{aligned}
k_- &= \left| \frac{1}{2} \left(p_0 \sqrt{1 - \frac{4m_q^2}{M^2}} - p \right) \right|, \\
k_+ &= \frac{1}{2} \left(p_0 \sqrt{1 - \frac{4m_q^2}{M^2}} + p \right). \qquad (7)
\end{aligned}
$$

The second term in the imaginary part, being proportional to $\theta(-M^2)$, describes Landau damping which takes place only for spacelike momenta, i.e. in scattering processes. Therefore it does not contribute to the dilepton production. In the hard thermal loop limit, $p_0, p \ll k$, the first term of ImF and ImG vanishes, whereas the second one reduces to the well known Landau damping contribution in this limit [9,20].

We do not take into account the vacuum part of the ρ self energy because the decay of a ρ-meson into a free quark-antiquark pair does not take place in the vacuum. The Lagrangian (1) can be considered as an effective Lagrangian valid only at temperatures at or above the phase transition.

Combining (3) to (6) we obtain the dilepton production rate (ρ-quark rate) by numerical integration over the loop momentum k (magnitude of the three-momentum). In Fig. 1 this rate is compared to the Born rate (direct quark-antiquark annihilation) [4] and the π^+-π^- annihilation rate via VMD (ρ-π rate) calculated by Gale and Kapusta [19]. In contrast to the Born rate the ρ-quark rate as well as the ρ-π rate show a clear peak in the vicinity of the ρ-mass. For $M < 0.4$ GeV the ρ-quark rate agrees well with the Born rate, which both have no mass cut assuming vanishing u- and d-quark masses. For $M > 0.4$ GeV

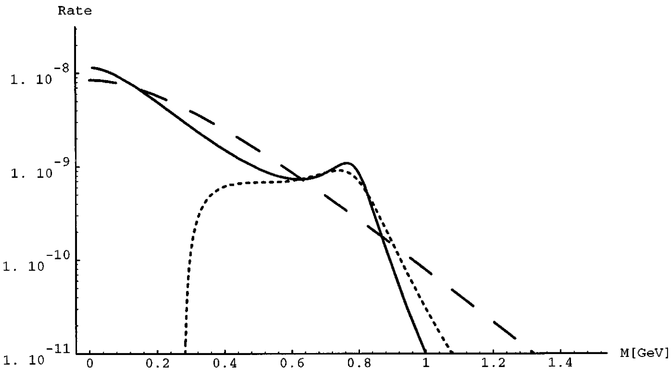

Fig. 1. ρ-quark rate (solid line), Born rate (dashed line), and ρ-π rate (dotted line) at $T = 0.15$ GeV, $p = 0.2$ GeV and $m_q = 0$

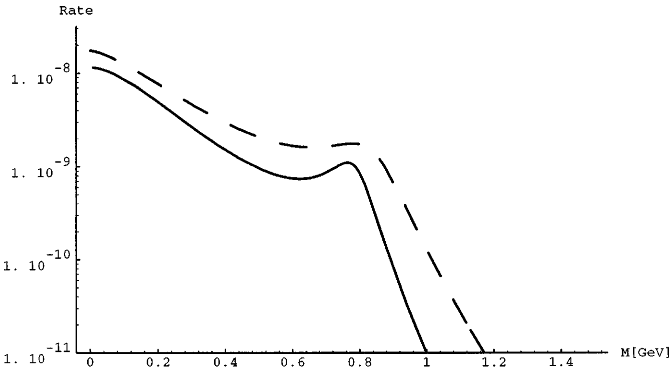

Fig. 2. ρ-quark rate at $T = 0.15$ GeV (solid line), $T = 0.2$ GeV (dashed line), $p = 0.2$ GeV, and $m_q = 0$

the ρ-quark rate agrees better with the ρ-π rate, which is zero for $M < 0.28$ GeV due to the finite π-mass. The agreement of the absolute values of these rates might be accidental as, for example, the ρ-quark rate is enhanced compared to the ρ-π rate by the larger number of quark degrees of freedom than π degrees of freedom, but reduced by the neglect of vacuum contributions in the first one.

We have chosen a temperature of $T = 0.15$ GeV in accordance with the critical temperature for the QCD phase transition as predicted by QCD lattice theory, since we expect ρ-mesons to co-exist with quarks only close to T_c. For the momentum of the virtual photon we have chosen $p = 0.2$ GeV, because CERES-SPS data show that the dilepton enhancement occurs at small transverse momenta 0.2 GeV$< p_T < 0.5$ GeV [1]. In Fig. 2 the ρ-quark rate is shown for different temperatures ($T = 0.15$ GeV and $T = 0.2$ GeV). In Fig. 3 the p-dependence of this rate is presented by comparing the rate at $p = 0.2$ GeV and $p = 1$ GeV.

Finally, we discuss briefly the possibility of an effective, temperature dependent quark mass as indicated by comparing the equation of state of an ideal gas of massive quarks and gluons with lattice calculations. In the temperature regime between T_c and about $2T_c$, quark masses of the order of $m = 500$ GeV are required to match the quasiparticle equation of state to lattice results [8]. In Fig. 4 the

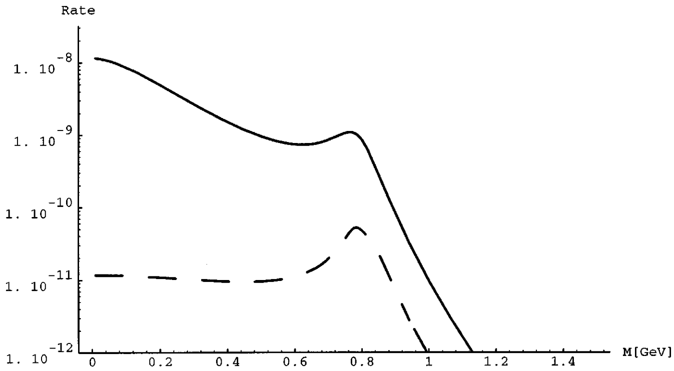

Fig. 3. ρ-quark rate at $T = 0.15$ GeV, $p = 0.2$ GeV (solid line), $p = 1$ GeV (dashed line), and $m_q = 0$

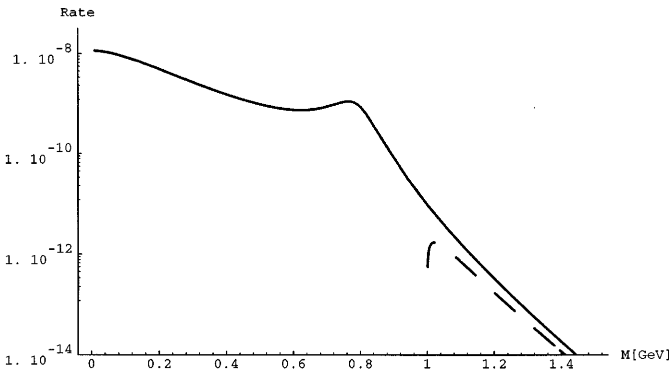

Fig. 5. ρ-quark rate at $T = 0.15$ GeV, $p = 0.2$ GeV, $m_q = 0$ (solid line), and $m_q = 0.5$ GeV (dashed line)

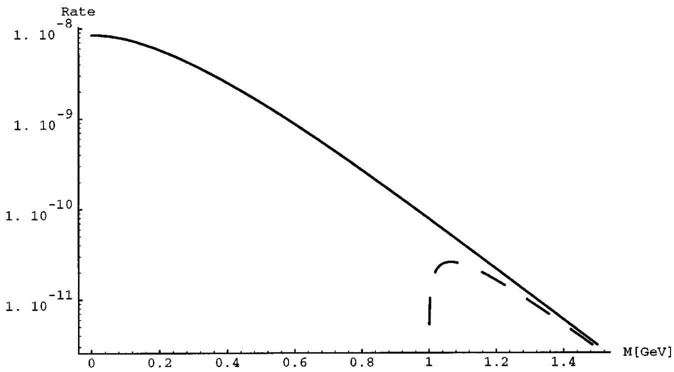

Fig. 4. Born rate at $T = 0.15$ GeV, $p = 0.2$ GeV, $m_q = 0$ (solid line), and $m_q = 0.5$ GeV (dashed line)

influence of such an effective quark mass on the Born rate is shown. The rate vanishes for $M < 2m_q$ but approaches quickly the bare mass rate above $M = 1$ GeV. This result has also been found by Peshier et al. [8], to which we would like to compare now our result of the ρ-quark rate (Fig. 5), where the ρ -peak now is absent in the case of finite quark masses.

If this simple picture for the QGP were true, there would be no dilepton production from the QGP below the ρ -peak, where the SPS dilepton enhancement has been observed. On the other hand, although the simple assumption of an effective quark mass might be sufficient to explain the equation of state, the correct quark dispersion relation in a QGP may look completely different. Perturbative as well as non-perturbative approaches indicate two quark branches starting from the same effective mass at zero momentum. The lower branch, corresponding to a collective quark mode (plasmino) that is absent in vacuum, shows a minimum at finite momentum. The splitting of the two collective quark modes and the minimum of the plasmino branch give rise to interesting structures, namely peaks (van Hove singularities) and gaps, in the dilepton rate. In particular, low mass dileptons are possible coming from electromagnetic transitions from the upper to the lower branch [6,13]. Furthermore, a finite width of the quarks from scattering in the QGP will change the dilepton production rate in addition.

Summarizing, we have calculated the dilepton production rate from quark-antiquark annihilation assuming the co-existence of quarks and ρ -mesons around the QCD phase transition as indicated by lattice QCD and gauged linear sigma model calculations. Using VMD the dilepton production rate has been calculated in a similar way as the rate from π^+-π^- rate, replacing the π-loop by a quark loop in the latter. For low invariant masses the rate turns out to agree well with the Born rate, whereas for higher invariant masses the agreement with the ρ-π rate is better. Although the surprisingly good coincidence in the absolute value of these rates may be accidental, considering for example possible uncertainites in the coupling constant g, the agreement of the shape is easy to understand (no mass cut for massless quarks, ρ -peak in VMD).

Concerning SPS data the contribution from the ρ-quark rate to the dilepton spectrum is by far too small to explain the observed enhancement [1]. For the ρ-quark rate is of the same order as the Born rate, which leads to a dilepton yield that is about two orders of magnitude smaller than the observed one according to hydrodynamical calculations [16]. However, this situation will be different at RHIC and LHC where a higher initial temperature and a longer lifetime of the QGP are expected.

Finally, we have shown that an effective quark mass as indicated by lattice calculations of the equation of state of the QGP leads to a suppression of dileptons with invariant masses below about 1 GeV. However, the simple picture of non-interacting massive quarks is probably too oversimplified for predicting the dilepton production from the QGP.

References

1. G. Agakichiev et al., CERES collaboration, Phys. Rev. Lett. **75**, 1272 (1995); N. Masera for the HELIOS-3 collaboration, Nucl. Phys. A **590**, 93c (1995); A. Drees for the CERES collaboration, Nucl. Phys. A **630**, 449c (1998); G. Agakichiev et al., CERES collaboration, Phys. Lett. B **422**, 405 (1998)
2. G.E. Brown and M. Rho, Phys. Rev. Lett. **66**, 2720 (1991); B. Friman and H.J. Pirner, Nucl. Phys. A **617**, 496 (1997); R. Rapp, G. Chanfray, and J. Wambach, Nucl. Phys. A

617, 472 (1997); F. Klingl, N. Kaiser, and W. Weise, Nucl. Phys. A **624**, 527 (1997); W. Peters, M. Post, H. Lenske, S. Leupold, and U. Mosel, Nucl. Phys. A **632**, 109 (1998); W. Cassing, E.L. Bratkovskaya, R. Rapp, and J. Wambach, Phys. Rev. C **57**, 916 (1998)
3. P.V. Ruuskanen, Nucl. Phys. A **544**, 169c (1992)
4. J. Cleymans, J. Fingberg, and K. Redlich, Phys. Rev. D **35**, 2153 (1987)
5. T. Altherr and P.V. Ruuskanen, Nucl Phys B **380**, 377 (1992); J. Cleymans and I. Dadić, Phys. Rev. D **47**, 160 (1993); J. Cleymans, I. Dadić, and J. Joubert, Phys. Rev. D **49**, 230 (1994)
6. E. Braaten, R.D. Pisarski, and T.C. Yuan, Phys. Rev. Lett. **64**, 2242 (1990); S.M.H. Wong, Z. Phys. C **53**, 465 (1992)
7. M.H. Thoma and C.T. Traxler, Phys. Rev. D **56**, 198 (1997)
8. A. Peshier, B. Kämpfer, O.P. Pavlenko, and G. Soff, Phys. Lett. B **337**, 235 (1994); A. Peshier, B. Kämpfer, O.P. Pavlenko, and G. Soff, Phys. Rev. D **54**, 2399 (1996)
9. E. Braaten and R.D. Pisarski, Nucl. Phys. B **337**, 569 (1990)
10. P. Aurenche, F. Gelis, R. Kobes, and H. Zaraket, Phys. Rev. D **58**, 085003 (1998); P. Aurenche, F. Gelis, R. Kobes, and H. Zaraket, hep-ph/9903307
11. G. Boyd et al., Nucl. Phys. B **469**, 419 (1996)
12. A. Schäfer and M.H. Thoma, Phys. Lett. B **451**, 195 (1999)
13. M.G. Mustafa, A. Schäfer and M.H. Thoma, in preparation
14. G. Boyd, S. Gupta, F. Karsch, and E. Laermann, Z. Phys. C **64**, 331 (1994)
15. R.D. Pisarski, Nucl. Phys. A **590**, 553c (1995); R.D. Pisarski, Phys. Rev. D **52**, 3773 (1996)
16. J. Sollfrank et al., Phys. Rev. C **55**, 392 (1997)
17. A. Polleri, R.A. Broglia, P.M. Pizzochero, and N.N. Scoccola, Z. Phys. A **357**, 325 (1997)
18. Y.B. He, J. Hüfner, S.P. Klevansky, and P. Rehberg, Nucl. Phys. A **630**, 719 (1998)
19. C. Gale and J.I. Kapusta, Nucl. Phys. B **357**, 65 (1991)
20. V.V. Klimov, Sov. Phys. JETP **55**, 199 (1982); H.A. Weldon, Phys. Rev. D **26**, 1394 (1982)
21. J.I. Kapusta, *Finite Temperature Field Theory* (Cambridge University Press, New York, 1989) p.71