Low-Mass Dileptons at the CERN-SpS: Evidence for Chiral Restoration?

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Using a rather complete description of the in-medium $\rho$ spectral function - being constrained by various independent experimental information - we calculate pertinent dilepton production rates from hot and dense hadronic matter. The strong broadening of the $\rho$ resonance entails a reminiscence to perturbative $q\bar{q}$ annihilation rates in the vicinity of the phase boundary. The application to dilepton observables in Pb(158 AGeV)+Au collisions - incorporating recent information on the hadro-chemical composition at CERN-SpS energies - essentially supports the broadening scenario. Possible implications for the nature of chiral symmetry restoration are outlined.

I. INTRODUCTION

In the low-energy sector ($q^2 \leq 1-2$ GeV$^2$) the properties of Quantum Chromodynamics are governed by (approximate) chiral symmetry and its dynamical breaking. In hot and dense matter as created through energetic collisions of heavy nuclei chiral symmetry is believed to be restored, associated with the disappearance of the chiral condensate and a substantial reshaping of the low-lying hadronic spectrum. As the chiral condensate does not represent a physical observable, the detection of signatures for chiral restoration has to focus on medium effects in hadronic properties. Dileptons as penetrating probes provide the opportunity to study rather directly the in-medium modifications of light vector mesons through their decays $V \to l^+l^-$.

Measurements of low-mass dilepton spectra in heavy-ion collisions at CERN-SpS energies\textsuperscript{11} have revealed appreciable radiation beyond final-state hadron decays. However, the spectral shape of this additional contribution - stemming from the interaction phase of the hadronic fireball - is not easily accounted for. In particular, using the vacuum vector meson line shapes, the enhancement around $M_V \simeq (0.2-0.6)$ GeV remains unexplained. At the same time the yield in the $\rho/\omega$ mass region tends to be overestimated. The central question then is whether one has possibly observed signatures of (the onset of) chiral symmetry restoration, and, if so, how it manifests itself in the vector channel.

Several theoretical approaches have been pursued to study medium effects in dilepton production rates from hot/dense matter, e.g., the ‘dropping rho mass’ scenario (within a mean-field type treatment)\textsuperscript{15,16}, the chiral reduction formalism (based on chiral Ward identities in connection with low-density expansions)\textsuperscript{17}, chiral Lagrangian frameworks\textsuperscript{18} or many-body-type calculations of in-medium vector-meson spectral functions\textsuperscript{18,19}. As a further complication, the quantitative impact of in-medium effects on the final spectra depends on another important ingredient, namely the modeling of the space-time evolution of the global heavy-ion collision dynamics. Here, transport calculations seem to give substantially larger dilepton yields than hydrodynamic simula-

II. PRODUCTION RATES IN HOT/DENSE MATTER

The radiation of dileptons from a hot and dense medium characterized by a temperature $T$ and baryon chemical potential $\mu_B$, \[ \frac{d^6 N_{l^+l^-}}{d^4\mathbf{q}^2} = -\frac{\alpha^2}{\pi^3M^2} f_B(q_0;T) \text{Im}\Pi_{\text{em}}(q_0,\mathbf{q};\mu_B,T), \] (1)

is governed by the thermal expectation value of the (retarded) electromagnetic current-current correlator, $\Pi_{\text{em}}(f_B; \text{Bose function})$. The latter can be written in terms of its longitudinal and transverse projections as \[ \Pi_{\text{em}}(q_0,\mathbf{q}) = \frac{1}{3} \left[ \Pi_{\text{em}}^L(q_0,\mathbf{q}) + 2\Pi_{\text{em}}^T(q_0,\mathbf{q}) \right], \] (2)

both depending separately on energy and momentum. Below momentum transfers of about 1 GeV, the hadronic part of the e.m. current can be accurately saturated by
the light vector mesons within the well-established vector dominance model (VDM). The correlator is then expressed through the imaginary parts of the vector meson propagators (= spectral functions) as

$$\text{Im}\Pi_{em} = \sum_{V=\rho,\omega,\phi} \frac{(m_V^{(0)})^4}{g_V^2} \text{Im} D_V.$$  \hspace{1cm} (3)

In the following we will neglect the small contributions from the isoscalar part and concentrate on the $\rho$ meson which plays by far the dominant role for the time scales involved in heavy-ion reactions.

Along the lines of our earlier analyses \cite{8,16} the $\rho$ propagator in hot hadronic matter,

$$D^{L,T}_{\rho} = \frac{1}{M^2 - (m_\rho^{(0)})^2 - \Sigma^{L,T}_{\rho\pi\pi} - \Sigma^{L,T}_{\rho M} - \Sigma^{L,T}_{\rho B}},$$  \hspace{1cm} (4)

is evaluated in terms of various contributions entering its in-medium selfenergy $\Sigma_\rho$. The meson gas effects (encoded in $\Sigma_{\rho M}$) are accounted for following Ref. \cite{15} through interactions with the most abundant thermal $\pi$, $K$ and $\rho$ mesons saturated with a rather complete set of $s$-channel resonances $R$ up to 1.3 GeV. The interaction vertices have been constrained by both hadronic ($R \to P\rho$) and radiative ($R \to P\gamma$) decay branchings. Also, the Bose enhancement for the in-medium $\rho\pi\pi$ width has been included in $\Sigma_{\rho\pi\pi}$.

The $\rho$ modifications in nuclear matter are based on the model constructed in Refs. \cite{8,14,17}, where direct $\rho N \to B$ interactions ($B=N$, $\Delta$, $N(1520)$, $\Delta(1700)$, $N(1720)$, ...), encoded in $\Sigma_{\rho B}$ as well as pion cloud modifications through $\pi NN^{-1}$ and $\pi\Delta N^{-1}$ excitations (entering $\Sigma_{\rho\pi\pi}$) were calculated. Again, the model parameters have been thoroughly constrained by analyzing photoabsorption spectra on nucleons and nuclei as well as $\pi N \to \rho N$ scattering data. The latter enforce a rather soft form factor on the $\pi NN$ vertex. Also, since the VDM is less accurate in the baryonic sector, an improved VDM coupling \cite{8} has been employed for the (transverse) $\gamma N \to B$ transition form factors so that

$$\text{Im}\Pi_{em} = \frac{1}{3g_{\rho\pi\pi}^2} \left[ F^L + 2F^T \right]$$  \hspace{1cm} (5)

with

$$F^L = -(m_\rho^{(0)})^4 \text{Im} D^L_\rho,$$

$$F^T = -\text{Im} \Sigma^{L,T}_{\rho\pi\pi} |d_\rho - 1|^2 - \text{Im} \Sigma^{L,T}_{\rho M} |d_\rho - r_B|^2,$$

$$d_\rho = \frac{M^2 - \Sigma^{L,T}_{\rho\pi\pi} + \Sigma^{L,T}_{\rho M} - r_B \Sigma^{L,T}_{\rho B}}{M^2 - (m_\rho^{(0)})^2 - \Sigma^{L,T}_{\rho\pi\pi} - \Sigma^{L,T}_{\rho M} - \Sigma^{L,T}_{\rho B}},$$  \hspace{1cm} (6)

cf. Refs. \cite{8,16} (here $r_B = 0.7$, the ratio between the actual $\gamma NB$ coupling and its value in the simple VDM).

At finite temperatures two additional features appear in the baryonic sector. Firstly, the Fermi-Dirac distribution functions are replaced by thermal ones \cite{18}, which, at the temperatures of interest here ($T \approx 150$ MeV), are substantially smeared as compared to the zero-temperature $\Theta$-functions. Secondly, a large fraction of the nucleons is excited into baryonic resonances which, in turn, are no longer active in nucleon-driven effects. However, as conjectured in Ref. \cite{18}, a baryonic resonance $B_1$ might still have excitations of type $\rho B_2 B_1^{-1}$ built on it. The particle data table \cite{19} indeed supports this hypothesis: e.g., the $\Sigma(1670)$ (which is a well-established four-star resonance with spin-isospin $IJK^P = 1^-_2$), when interpreted as a $\rho \Sigma$ (or $\rho A$) ‘resonance’, very much resembles the quantum numbers and excitation energy ($\Delta E \approx 500 - 700$ MeV) of the $\rho N \to N(1520)$ transition. In addition, the branching ratio of $\Sigma(1670)$ decays into ‘simple’ final states such as $NK$, $\Sigma\pi$ or $\Lambda\pi$ is significantly less than 100%. Similar excitations on non-strange baryonic resonances are more difficult to identify as the latter themselves decay strongly via pion emission (i.e., the $B_2 \to B_1 \rho$ decay is immediately followed by further $B_2 \to \pi N$ and $\rho \to \pi\pi$ decays). Neverthe-
less, from pure quantum numbers it is tempting to associate, e.g., \( \Delta(1930) \Delta^{-1} \) or \( N(2080)N(1440)^{-1} \) excitations with \( S \)-wave 'Rhosobar' states. We have conservatively estimated their coupling constants to give resonances in the vicinity of heavy-quark thresholds such as \( \rho_B \) and \( \gamma_B \) (the latter are typically in the range of 0.2–0.7 MeV). Due to the large total widths involved, their impact on the \( \rho \) propagator is weaker as compared to excitations on nucleons at equal densities; e.g., at \( T=150 \) MeV and \( \mu_B = 436 \) MeV, where, in chemical equilibrium, about 2/3 of the baryons are in excited states, the additional broadening from the calculated \( B_2 B_1^{-1} \) excitations as compared to \( BN^{-1} \) ones amounts to about 40%. Concerning the pion cloud modifications, we do not explicitly compute the 'Pisobars' on excited resonances but approximate their effect by using an effective nucleon density \( q_{eff} = \rho_N + 0.5 \rho_B^* \).

The final result for the \( \rho \) spectral function in hot hadronic matter is displayed in Fig. 1 at fixed three-momentum and chemical potentials \( \mu_\pi = 0, \mu_B = 330 \) MeV for pions and baryons, respectively. One observes a strong broadening with increasing temperature and density. Comparing to the pure meson gas results of Ref. [15] (cf. Fig. 3 therein), it becomes clear that especially the low-mass enhancement around \( M \approx 0.4 \) GeV is largely driven by baryons.

The pertinent three-momentum integrated dilepton production rates are shown in Fig. 2. At moderate temperature and density (upper panel), one still recognizes a remnant of the \( \rho \) peak, which, however, is entirely wiped out under conditions expected to be close to the phase boundary (lower panel). This provokes a comparison with quark-gluon based rate calculations, which, for simplicity, have been evaluated in terms of lowest order \( O(\alpha_s^0) \) \( q\bar{q} \) annihilation [22],

\[
\frac{dR_{q\bar{q}\to ee}}{d^2q} = \frac{\alpha_s^2}{4\pi^4} \frac{T}{f_B(q_0; T)} \sum_q e_q^2 \frac{(x_- + y) (x_+ + \exp[-\mu_q/T])}{(x_- + y) (x_+ + \exp[-\mu_q/T])} \tag{7}
\]

with \( x_\pm = \exp[-(q_0 \pm q)/2T], y = \exp[-(q_0 + \mu_q)/T] \). In the vacuum the \( q\bar{q} \) rates are known to coincide with the hadronic description for invariant masses \( M \geq 1.5 \) GeV as marked by the cross section ratio \( \sigma(e^+e^- \to \text{hadrons})/\sigma(e^+e^- \to \mu^+\mu^-) \) for the inverse process of \( e^+e^- \) annihilation (up to additional resonance structures in the vicinity of heavy-quark thresholds such as \( c\bar{c} \)). It seems that the in-medium hadronic rates indeed approach the partonic ones rather quickly leading to an approximate agreement at the highest temperatures/densities for masses of about 0.5–1 GeV (the deviation at masses \( M > 1 \) GeV is due to the incomplete description of the vector correlator in vacuum which does not include more than two-pion states; the discrepancy at low masses might be reduced once higher order \( \alpha_s \) corrections are included, in particular soft Bremsstrahlung-type processes). A tempting interpretation of this behavior is a lowering of the 'quark-hadron duality threshold' in hot and dense hadronic matter. We will return to this issue in the discussion in Sect. IV.

\[\text{III. LOW-MASS DILEPTON OBSERVABLES IN PB(158 AGeV)+AU}\]

For a sensible application of the in-medium vector correlator to calculate dilepton production in URHIC’s a realistic temperature and density profile is required. Microscopic transport or hydrodynamical simulations, e.g., have proven very successful in describing the final hadron...
observables. However, the associated dilepton yields, based on identical production rates, may differ appreciably, cf., e.g., Refs. [24] and [25]. As we will show below, a possible origin of this discrepancy might be related to the build-up of a finite chemical potential, which is usually not included in most hydrodynamical calculations to date (see, however, Ref. [26]).

Recent hydro-chemical analyses [12,13] of a large body of hadronic heavy-ion data has shown that the finally observed particle abundances at SpS energies are consistent with a common chemical freezeout at temperatures/baryon chemical potentials around (T, µ_B)ch = (175, 270) MeV. In the subsequent expansion and cooling the system is still strongly interacting with elastic collisions maintaining thermal equilibrium until the \textit{thermal freezeout}. The absence of pion-number changing (inelastic) reactions (which is well supported phenomenologically from the inelasticities in \(\pi\pi\) scattering) then induces a finite pion chemical potential [4]. To incorporate these features, we here take recourse to a simple expanding thermal fireball model. We start by generating a (non-interacting) resonance gas equation of state of hot hadronic matter including the lowest 32 (16) mesonic (baryonic) states. Imposing entropy and baryon number conservation determines the temperature dependence of the baryon chemical potential, \(\mu_B(T)\). The additional requirement of pion- [kaon-] number conservation entails a finite \(\mu_\pi(T)\) \((\mu_K(T) \simeq \mu_K(T))\), with the various baryon chemical potentials kept in relative equilibrium (with respect to strong interactions), e.g., \(\mu_\Delta(T) = \mu_N(T) + \mu_\pi(T)\), etc.. With an entropy per baryon of \(s/\rho_B = 26\), our trajectory is found to pass through (T, µ_N) = (175, 250) MeV, compatible with the experimentally deduced point [13] (here, a small \(\mu_\pi \simeq 20\) MeV is needed to be consistent with the final pion-to-baryon ratio of about 5:1 at SpS energies). Towards thermal freezeout at (T, \(\mu_N\))fo \(\simeq (110 - 120, 415 - 450)\) MeV, \(\mu_\pi(T)\) and \(\mu_K(T)\) increase approximately linearly to around 80 and 110 MeV, respectively. Finally we need to introduce a time scale to obtain the volume expansion. We approximate the latter by a cylindrical geometry as

\[
V_{FC}^{(2)}(t) = 2 \left( z_0 + v_z t + \frac{1}{2} a_z t^2 \right) \pi \left( r_0 + \frac{1}{2} a_\perp t^2 \right)^2
\]  

(8)

employing two firecylinders expanding in the \(\pm z\) direction. Guided by hydrodynamical simulations [26] the primordial longitudinal motion for Pb(158 AGeV)+Au reactions is taken to be \(v_z = 0.5c\), and the longitudinal and transverse acceleration are fixed to give final velocities \(v_z(t_{fo}) \simeq 0.75c\), \(v_\perp(t_{fo}) \simeq 0.55c\) as borne out from experiment (this, in turn, requires fireball lifetimes of about \(t_{fo} = 10 - 12\) fm/c and implies transverse expansion by 3-4 fm, consistent with HBT analyses [27]). The parameter \(r_0\) denotes the initial nuclear overlap radius, e.g., \(r_0 = 4.6\) fm for collisions with impact parameter \(b = 5\) fm and \(N_p \approx 260\) participant baryons. The parameter \(z_0\) is equivalent to a formation time and fixes the starting point of the trajectory in the (\(T, \mu_N\)) plane.

With the such specified space-time evolution for \(\sim 30\%\) central Pb(158 AGeV)+Au collisions the dilepton yields from in-medium radiation are obtained as

\[
d^2 N_{ee} / dq dM = \frac{\alpha^2}{\pi^3 g_{\rho \pi \pi} M} \int_0^t \frac{dt}{\tau_c} V_{FC}^{(1)}(t) \int \frac{d^3q}{q_0} f^0(q_0; T(t))
\]

\[
\times F(M, q; \mu_B(t), T(t)) \ Acc(M, \vec{q})
\]  

(9)

with Acc\((M, \vec{q})\) accounting for the experimental acceptance cuts (as well as an approximate mass resolution) in the ‘96 CERES data. The backward firecylinder is placed at a rapidity \(y = 2.6\) to ensure the correct charged particle multiplicity as quoted by CERES. The electromagnetic ‘transition form factor’ \(F\) contains the in-medium \(\rho\) spectral function as described above, where the additional effects of the finite meson chemical potentials have been implemented in Boltzmann approximation. For initial conditions \((T, \rho_B)_{ini} = (190 \ MeV, 2.55 \rho_0)\) and a freeze-out at \((T, \rho_B)_{fo} = (115 \ MeV, 0.33 \rho_0)\) the fireball lifetime amounts to about \(t_{fo} = 11\) fm/c, see also above (reducing, e.g., \((T, \rho_B)_{ini}\) to \((180 \ MeV, 1.92 \rho_0)\) reduces the dilepton signal from the fireball by about 15%).

\begin{figure}
\centering
\includegraphics[width=\textwidth]{fig3.png}
\caption{1996 CERES/NA45 inclusive dielectron invariant mass spectra from 30\% central Pb(158 AGeV)+Au collisions [26] compared to the hadronic cocktail [28] (without \(\rho\) decays, dashed-dotted line) and to the cocktail plus \(\pi\pi\) annihilation from the hadronic fireball using either the free (dashed line) or the in-medium (solid line) \(\rho\) spectral function from Eq. [4].}
\end{figure}

In addition to the in-medium radiation described by Eq. [4], the finally observed dilepton spectra contain a sizeable contribution from free hadron decays after freezeout, the so-called ‘cocktail’. For this we employ the most recent evaluation of the CERES collaboration [29], which is generated from particle abundances based on the chemical freezeout of Ref. [13] being consistent with our \((T, \mu_B)\) trajectory. As in our previous works [8] the \(\rho\)-meson decays have been excluded from the cocktail as,

\[
\begin{align*}
\end{align*}
\]
owing to the short lifetime of the $\rho$, they are accounted for through the in-medium part, Eq. (1), in the vicinity of the (idealized) 'freezeout'.

The final results for the inclusive invariant mass spectra, shown in Fig. 3, demonstrate that the use of the in-medium spectral function leads to reasonable agreement with the ’96 CERES data, which cannot be described assuming vacuum $\rho$ properties. This is in line with the conclusions of our earlier analysis [8], but we emphasize again that the results presented here are based on a much improved understanding of both the microscopic $\rho$ spectral function (being constrained by a large body of independent data) as well as the space-time evolution dynamics (consistent with chemical freezeout analyses, without the use of an overall normalization factor, etc.). As compared to Ref. [8] the smaller $\mu_B$ in the early stages, the finite $\mu_\pi$ in the later stages, and the more complete assessment of the meson gas effects [15] reduces the dominant role of the baryonic contributions to about equal importance as the mesonic ones. Another noteworthy point is that the overall dilepton yield exceeds what has been found in the hydrodynamical calculations of, e.g., Ref. [25] (where $\mu_\pi = 0$ has been used throughout) by about a factor of $\sim 2.5$, thereby essentially resolving the above mentioned discrepancy to transport calculations [15].

A more detailed insight into the nature of the low-mass enhancement is revealed by inspecting its dependence on the transverse pair momentum $q_t$. Fig. 4 shows the invariant mass spectrum separated into two transverse momentum bins below and above $q_t = 0.5$ GeV. Clearly, the excess of the data over the free calculation is most pronounced in the small momentum bin which is properly reflected in our in-medium calculations. The same behavior is exhibited in another projection of the data in Fig. 5, where transverse momentum spectra in various invariant mass bins are displayed.

![FIG. 5. 1996 CERES/NA45 transverse momentum spectra for $e^+e^-$ pairs in 30% central Pb(158 AGeV)+Au collisions 28. Line identification as in Fig. 3.](image)

**IV. SUMMARY AND DISCUSSION**

Based on realistic hadronic interaction vertices we have shown that the in-medium $\rho$ spectral function, evaluated within standard many-body techniques, undergoes a strong broadening in hot and dense matter. At comparable densities the baryonic effects prevail over mesonic ones due to the stronger nature of the meson-nucleon interactions. Corresponding thermal dilepton production rates exhibit a remarkable tendency to approach the lowest-order perturbative $q\bar{q}$ annihilation with increasing temperature and density. This might indicate the emergence of 'quark-hadron duality' towards the phase boundary down to rather low mass scales of 0.5 GeV, made possible through large imaginary parts resummed in the Dyson equation of $\rho$ propagator. The medium effects might thus be interpreted as a lowering of the 'duality threshold', which in free space is located around 1.5 GeV. An unambiguous identification of chiral restoration, however, requires the simultaneous treatment of the axial correlator ($a_1$ channel) which has to merge with the vector correlator which we have focused on here. Although the perturbative $q\bar{q}$ rate implies restored chiral
symmetry, one has to keep in mind that (thermal) non-perturbative effects below 1 GeV are not yet well under control. The argument can be made more rigorous in the 1–1.5 GeV region, where the lowest-order in temperature mixing between vector and axialvector correlators suffices to establish a full degeneracy between them as well as the perturbatively calculated partonic result.

We have furthermore demonstrated that our model for the in-medium $\rho$ spectral function is in line with current low-mass dilepton measurements at the CERN-SpS. For that we have employed a simple fireball expansion model consistent with recent hadro-chemical analysis and imposing effective pion-number conservation between chemical and thermal freezeout via finite pion chemical potentials (the latter seem to resolve a large part of the discrepancy in the total dilepton yield between transport and chemical equilibrium-based hydrodynamic simulations). The apparent depletion of the in-medium dilepton yield in the $\rho/\omega$ region, as well as the enhancement below can be accounted for. This also holds for the low-$q_t$ nature of the excess. Upcoming high resolution measurements as well as the more baryon-dominated 40 GeV run will put the model predictions under further scrutiny and can be expected to advance our understanding of chiral symmetry restoration in hot and dense hadronic matter.

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[1] G. Agakichiev et al., CERES collaboration, Phys. Rev. Lett. 75 (1995) 1272; P. Wurm for the CERES collaboration, Nucl. Phys. A590 (1995) 103c;
[2] N. Masera for the HELIOS-3 collaboration, Nucl. Phys. A590 (1995) 93c;
[3] G.E. Brown and M. Rho, Phys. Rev. Lett. 66 (1991) 2720.
[4] W. Cassing, W. Ehehalt and C.M. Ko, Phys. Lett. B363 (1995) 35.
[5] G.Q. Li, C.M. Ko and G.E. Brown, Phys. Rev. Lett. 75 (1995) 4007; G.Q. Li, C.M. Ko, H. Sorge and G.E. Brown, Nucl. Phys. A611 (1996) 539.
[6] J.V. Steele, H. Yamagishi and I. Zahed, Phys. Lett. B384 (1996) 255; Phys. Rev. D56 (1997) 5605.
[7] C. Song and V. Koch, Phys. Rev. C54 (1996) 3218.
[8] R. Rapp, G. Chanfray and J. Wambach, Nucl. Phys. A617 (1997) 472; G. Chanfray, R. Rapp and J. Wambach, Phys. Rev. Lett. 76 (1996) 368.
[9] B. Friman and H.J. Pirner, Nucl. Phys. A617 (1997) 496.
[10] F. Klingl, N. Kaiser and W. Weise, Nucl. Phys. A624 (1997) 527.
[11] W. Peters, M. Post, H. Lenske, S. Leupold and U. Mosel, Nucl. Phys. A632 (1998) 109.
[12] P. Braun-Munzinger and J. Stachel, Nucl. Phys. A638 (1998) 3c; P. Braun-Munzinger, J. Stachel, J.P. Wessels and N. Xu, Phys. Lett. B365 (1996) 1.
[13] J. Cleymans and K. Redlich, hep-th/9903063.
[14] H. Bebie, P. Gerber, J.L. Goity and H. Leutwyler, Nucl. Phys. B378 (1992) 95.
[15] R. Rapp and C. Gale, Phys. Rev. C (1999) in press, and hep-ph/9902268.
[16] R. Rapp, M. Urban, M. Buballa and J. Wambach, Phys. Lett. B417 (1998) 1.
[17] M. Urban, M. Buballa, R. Rapp and J. Wambach, Nucl. Phys. A641 (1998) 433.
[18] N.M. Kroll, T.D. Lee and B. Zumino, Phys. Rev. 157 (1967) 1376.
[19] R. Rapp and J. Wambach, Nucl. Phys. A573 (1994) 626.
[20] G.E. Brown, G.Q. Li, R. Rapp, M. Rho and J. Wambach, Acta Phys. Polon. B29 (1998) 2309.
[21] Particle Data Group, R.M. Barnett et al., Phys. Rev. D54 (1996) 1.
[22] J. Cleymans, J. Fingberg and K. Redlich, Phys. Rev. D35 (1987) 2153.
[23] R. Rapp, Proc. of Quark Matter ’99, Torino (Italy), May 10-15, 1999, and hep-ph/9907342.
[24] W. Cassing, E.L. Bratkovskaya, R. Rapp and J. Wambach, Phys. Rev. C57 (1998) 916.
[25] P. Huovinen and M. Prakash, Phys. Lett. B450 (1999) 15.
[26] C.M. Hung and E. Shuryak, Phys. Rev. C57 (1998) 1891.
[27] U.A. Wiedemann, Proc. of Quark Matter ’99, Torino (Italy), May 10-15, 1999, and nucl-th/9907048.
[28] G. Agakichiev et al., CERES collaboration, Phys. Lett. B422 (1998) 405; B. Lenkeit, Doctoral Thesis, University of Heidelberg, 1998.
[29] D. Irmscher, Doctoral Dissertation, Univ. Heidelberg 1993; T. Ullrich, Doctoral Dissertation, Univ. Heidelberg 1994; G. Agakichiev et al., CERES collaboration, Eur. Phys. J. C4 (1998) 231; and A. Drees, private communication.