Hypermatter in chiral field theory

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Abstract. A generalized Lagrangian for the description of hadronic matter based on the linear $SU(3)_L \times SU(3)_R$ σ-model is proposed. Besides the baryon octet, the spin-0 and spin-1 nonets, a gluon condensate associated with broken scale invariance is incorporated. The observed values for the vacuum masses of the baryons and mesons are reproduced. In mean-field approximation, vector and scalar interactions yield a saturating nuclear equation of state. Finite nuclei can be reasonably described, too. The condensates and the effective baryon masses at finite baryon density and temperature are discussed.

1. Introduction

Although the underlying theory of strong interactions is known, there is presently little hope to solve QCD at nonzero baryon density.
Instead, we formulate an effective theory based on symmetries which hopefully reflects the basic features of QCD in a solvable manner. Chiral and flavor symmetries, as well as scale invariance are used as guidelines to construct and to explore a σ-model type $SU(3)_L \times SU(3)_R$ Lagrangian at finite strangeness, baryon density and temprature.

2. Theory

Numerous attempts exist to apply the idea of chiral symmetry to nuclear physics. A convenient approach is to take the linear σ model introduced by Gell-Mann and Levy. However, wenn calculating the equation of state for nuclear matter with this model, bifurcations and multiple solutions lead to an unphysical equation of state. This suggests that an important physics ingredient is missing. We show that the inclusion of (broken) scale invariance allows for a satisfactory description of the hadronic mass spectrum, nuclear matter and finite nuclei. How this symmetry is incorporated in an effective Lagrangian at mean-field level will be discussed in subsection 2.3.

2.1. Nonlinear Realisation of Chiral Symmetry

We construct a chiral Lagrangian adopting a nonlinear realization of chiral symmetry as discussed in [1, 2, 3]. The idea is to use a representation where the heavy particles (i.e. baryons and scalar mesons) transform equally under left and right rotations. To accomplish this, it is necessary to dress these particles nonlinearly with pseudoscalar mesons. The application of this method to our approach has the advantage that the interaction between heavy particles is only governed by $SU(3)_V$. In addition, pseudoscalar particles appear explicitly in the Lagrangian only if derivative or explicit symmetry breaking terms are included. In the following we will outline the argumentation. For a thorough discussion, see [4].
Let us write the elementary spinors (=quarks) q which transform in the whole chiral space $SU(3)_L \times SU(3)_R$ into 'new' quarks $\tilde{q}$ by

$$q_L(x) = U(x)\tilde{q}_L(x) \qquad q_R(x) = U^\dagger(x)\tilde{q}_R(x) \tag{1}$$

with the pseudoscalar octet π_a arranged in $U(x) = \exp[-i\pi_a\lambda^a/2]$. From the algebraic composition of mesons in terms of quarks (e.g., spin-0 mesons have the form $M = \Sigma + i\Pi \simeq \overline{q}_L q_R + \overline{q}_L\gamma_5 q_R$), it is straightforward to transform from 'old' mesons Σ and Π (which transform in the whole $SU(3)_L \times SU(3)_R$ space) into 'new' mesons X and Y (which transform in $SU(3)_{L+R} = SU(3)_V$ space):

$$M = \Sigma + i\Pi = U(X + iY)U \quad . \tag{2}$$

Here, the parity-even part X is associated with the scalar nonet, whereas Y is taken to be the pseudoscalar singlet [5]. All hadron multiplets are expanded in a basis of Gell-Mann matrices, e.g., $X = \frac{1}{\sqrt{2}}\sum_{a=0}^{8}\xi_a\lambda_a$.

2.2. Baryon-meson interaction

In a similar way, the 'old' baryon octet Ψ forming the representation (8,1) and (1,8) is transformed into a 'new' baryon octet $B = \frac{1}{\sqrt{2}}\sum_{a=1}^{8} b_a\lambda_a$:

$$\Psi_L = UB_LU^\dagger \qquad \Psi_R = U^\dagger B_R U \quad . \tag{3}$$

The transformations of the exponential U are known [2, 3],

$$U' = LUV^\dagger = VUR^\dagger, \tag{4}$$

and by using the transformation properties of the 'old' fields (which can be deduced once the algebraic composition in terms of the chiral quarks is known) the 'new' baryons B and the 'new' scalar mesons X can be shown to transform in $SU(3)_V$.
The pseudoscalars reappear in the transformed model as the parameters of the symmetry transformation. Therefore, chiral invariants (without space-time derivatives) are

independent of the Goldstone bosons. The 'new' fields allow for invariants which are forbidden for the 'old' fields by chiral symmetry: The invariant linear interaction terms of baryons with scalar mesons are

$$\mathcal{L}_{BX} = g_8^S \mathrm{Tr}\left(\alpha[\overline{B}BX]_F + (1-\alpha)[\overline{B}BX]_D\right) + g_1^S \mathrm{Tr}\cdot(\overline{B}B)\mathrm{Tr}X\,, \qquad (5)$$

with antisymmetric $[\overline{B}BX]_F := \mathrm{Tr}(\overline{B}BX - \overline{B}XB)$ and symmetric $[\overline{B}BX]_D := \mathrm{Tr}(\overline{B}BX + \overline{B}XB) - \frac{2}{3}\mathrm{Tr}(\overline{B}B)\mathrm{Tr}M$ invariants. The masses of the whole baryon multiplet are generated spontaneously by the vacuum expectation values (VEV) of only *two* meson condensates: From the spin-0 fields only the VEV of the components proportional to λ_0 and the hypercharge $Y \sim \lambda_8$ are nonvanishing, and the vacuum expectation value $\langle M\rangle$ reduces to:

$$\langle M\rangle \frac{1}{\sqrt{2}}(\xi_0\lambda_0 + \xi_8\lambda_8) \equiv \mathrm{diag}\ (\frac{\sigma}{\sqrt{2}}, \frac{\sigma}{\sqrt{2}}, \zeta) \quad ,$$

in order to preserve parity invariance and assuming, for simplicity, $SU(2)$ symmetry of the vacuum. Hence, the baryon masses read:

$$\begin{aligned} m_N &= m_0 - \frac{1}{3}g_8^S(4\alpha-1)(\sqrt{2}\zeta - \sigma) \qquad (6)\\ m_\Lambda &= m_0 - \frac{2}{3}g_8^S(\alpha-1)(\sqrt{2}\zeta - \sigma)\\ m_\Sigma &= m_0 + \frac{2}{3}g_8^S(\alpha-1)(\sqrt{2}\zeta - \sigma)\\ m_\Xi &= m_0 + \frac{1}{3}g_8^S(2\alpha+1)(\sqrt{2}\zeta - \sigma) \end{aligned}$$

with $m_0 = g_1^S(\sqrt{2}\sigma + \zeta)/\sqrt{3}$. The three parameters g_1^S, g_8^S and α can be used to fit the baryon masses to their experimental values (table 1). Then, no additional explicit symmetry breaking term is needed. For $\zeta = \sigma/\sqrt{2}$ (i.e. $f_\pi = f_K$), the masses are degenerate, and the vacuum is $SU(3)_V$ invariant.

Although the construction of invariants is only governed by $SU(3)_V$, relations following from chiral symmetry as PCAC and the Goldberger-Treiman relation are incorporated. The model also allows to predict the masses of the meson nonet at zero and finite density. The interaction of the vector meson nonet $V_\mu = \frac{1}{\sqrt{2}}\sum_{i=0}^{8} v_\mu^i \lambda_i$ and the axial vector meson nonet $A_\mu = \frac{1}{\sqrt{2}}\sum_{i=0}^{8} a_\mu^i \lambda_i$ with baryons reads:

$$\mathcal{L}_{BV} = g_8^V \mathrm{Tr}(\overline{B}\gamma^\mu[V_\mu, B] + \overline{B}\gamma^\mu\{A_\mu\gamma_5, B\}) + g_1^V \mathrm{Tr}(\overline{B}B)\gamma^\mu \mathrm{Tr}(V_\mu + A_\mu) . (7)$$

In the mean-field treatment, the axial mesons have a zero VEV. The relevant fields in the SU(2) invariant vacuum, v_μ^0 and v_μ^8, are taken to have the ideal mixing angle $\sin\theta_v = \frac{1}{\sqrt{3}}$, yielding

$$\begin{aligned} \phi_\mu &= v_\mu^8 \cos\theta_v - v_\mu^0 \sin\theta_v = \frac{1}{\sqrt{3}}(\sqrt{2}v_\mu^0 + v_8^\mu) \qquad (8)\\ \omega_\mu &= v_\mu^8 \sin\theta_v + v_\mu^0 \cos\theta_v = \frac{1}{\sqrt{3}}(v_\mu^0 - \sqrt{2}v_8^\mu) \quad . \end{aligned}$$

For $g_1^V = g_8^V$, the strange vector field $\phi_\mu \sim \bar{s}\gamma_\mu s$ does not couple to the nucleon, and the coupling constant $g_{N\omega}$ is fixed to the binding energy of nuclera matter. The remaining couplings to the strange baryons are then determined by symmetry relations:

$$g_{\Lambda\omega} = g_{\Sigma\omega} = 2g_{\Xi\omega} = \frac{2}{3}g_{N\omega} = 2g_8^V \qquad g_{\Lambda\phi} = g_{\Sigma\phi} = \frac{g_{\Xi\phi}}{2} = \frac{\sqrt{2}}{3}g_{N\omega} \quad , \tag{9}$$

where their relative values are related to the additive quark model. For isospin asymmetric matter, the ρ-field has to be taken into account. Its coupling constants are also determined by SU(3) symmetry to be

$$g_{N\rho} = g_{\Xi\rho} = g_{\Sigma\rho}/2 = g_{N\omega}/3 \quad g_{\Lambda\rho} = 0. \tag{10}$$

2.3. Chirally invariant potential

The chirally invariant potential includes the mass terms for mesons, their self-interaction and the dilaton potential for the breaking of scale symmetry. For the spin-0 mesonic potential we take all independent combinations of mesonic self-interaction terms up to fourth order

$$\mathcal{V}_0 = \frac{1}{2}k_0\chi^2 \mathrm{Tr}M^\dagger M - k_1(\mathrm{Tr}M^\dagger M)^2 - k_2\mathrm{Tr}(M^\dagger M)^2 \tag{11}$$
$$-k_3\chi(\det M + \det M^\dagger) + k_4\chi^4 + \frac{1}{4}\chi^4 \ln\frac{\chi^4}{\chi_0^4} - \frac{\delta}{3}\chi^4 \ln\frac{\det M + \det M^\dagger}{2\det\langle M\rangle} \quad .$$

The constants k_i are fixed by the vacuum masses of the pseudoscalar and scalar mesons, respectively. These are determined by calculating the second derivative of the potential in the ground state.

The quadratic and cubic form of the interaction is made scale invariant by multiplying it with an appropriate power of the dilaton field χ. Originally, the dilaton field was introduced by Schechter in order to mimic the trace anomaly of QCD $\theta_\mu^\mu = \frac{\beta_{QCD}}{2g}\mathcal{G}^a_{\mu\nu}\mathcal{G}^{\mu\nu}_a$ in an effective Lagrangian at tree level [8]. The effect† of the logarithmic term $\sim \chi^4 \ln\chi$ is to break the scale invariance of the model Lagrangian so that the proportionality $\theta_\mu^\mu \sim \chi^4$ holds. The comparison of the trace anomaly of QCD with that of the effective theory allows for the identification of the χ-field with the gluon condensate:

$$\theta_\mu^\mu = \langle\frac{\beta_{QCD}}{2g}\mathcal{G}^a_{\mu\nu}\mathcal{G}^{\mu\nu}_a\rangle \equiv (1-\delta)\chi^4 \quad . \tag{12}$$

The parameter δ originates from the second logarithmic term with the chiral and parity invariant combination $\sim \det M + \det M^\dagger$. The term is a SU(3)-extension of the logarithmic term proportional to $\chi^4\ln(\sigma^2+\pi^2)$ introduced in [9]. An orientation for the value of δ may be taken from β_{QCD} at one loop level, which suggests the value

† According to [8], the argument of the logarithm has to be chirally and parity invariant. This is fulfilled by the dilaton which is a chiral singlet and a scalar.

$\delta = 2/11$ for three flavors and three colors. This value gives the order of magnitude about which the parameter δ will be varied.
For the spin-1 mesons a mass term is needed. The simplest, scale invariant form

$$\mathcal{L}^1_{vec} = \frac{1}{2} m_V^2 \frac{\chi^2}{\chi_0^2} \mathrm{Tr}(V_\mu V^\mu + A_\mu A^\mu) \tag{13}$$

(χ_0 is the VEV of the dilaton field) implies a mass degeneracy for the meson nonet. To split the masses one can add the chiral invariant [10]

$$\mathcal{L}^2_{vec} = \frac{1}{8} \mu \mathrm{Tr}[(F_{\mu\nu} + G_{\mu\nu})^2 M^\dagger M + (F_{\mu\nu} - G_{\mu\nu})^2 M^\dagger M] \quad , \tag{14}$$

where $F_{\mu\nu}$ and $F_{\mu\nu}$ are the field-strength tensors of the vector-and axial-vector fields, respectively. In combination with the kinetic energy term (Eq. 18), one obtains for the vector mesons

$$\begin{aligned} \mathcal{L}_{vec} = & -\frac{1}{4}[1-\mu\frac{\sigma^2}{2}](F_\rho^{\mu\nu})^2 - \frac{1}{4}[1-\frac{1}{2}\mu(\frac{\sigma^2}{2}+\zeta^2)](F_{K^*}^{\mu\nu})^2 \\ & -\frac{1}{4}[1-\mu\frac{\sigma^2}{2}](F_\omega^{\mu\nu})^2 - \frac{1}{4}[1-\mu\zeta^2](F_\phi^{\mu\nu})^2 \end{aligned} \tag{15}$$

Since the coefficients are no longer unity, the vector meson fields have to be renormalized, i.e., the new ω-field reads $\omega_r = Z_\omega^{-1/2}\omega$. The renormalization constants are the coefficients in the square brackets in front of the kinetic energy terms of Eq. (15), i.e., $Z_\omega^{-1} = 1 - \mu\sigma^2/2$. The mass terms of the vector mesons deviate from the mean mass m_V by the renormalization factor† , i.e.,

$$m_\omega^2 = m_\rho^2 = Z_\omega m_V^2 \quad ; \quad m_{K^*}^2 = Z_{K^*} m_V^2 \quad ; \quad m_\phi^2 = Z_\phi m_V^2 \quad . \tag{16}$$

The constant μ is fixed to give the correct ω-mass. The other vector meson masses are displayed in table 1.

2.4. Explicit breaking of chiral symmetry

The term

$$\mathcal{V}_{SB} = \frac{1}{2}\frac{\chi^2}{\chi_0^2} \mathrm{Tr} f(M + M^\dagger) \tag{17}$$

breaks the chiral symmetry explicitly and makes the pseudoscalar mesons massive. It is scaled appropriately to have dimensions equal to that of the quark mass term $\sim m_q \bar{q} q + m_s \bar{s} s$, which is present in the QCD Lagrangian with massive quarks. This term leads to a nonvanishing divergence of the axial currents. The matrix elements of $f = 1/\sqrt{2}(f_0\lambda_0 + f_8\lambda_8)$ can be written as a function of $m_\pi^2 f_\pi$ and $m_K^2 f_K$ to satisfy the (approximately valid) PCAC relations for the π- and K-mesons. Then, by utilizing the equations of motion, the VEV of σ and ζ are fixed in terms of f_π and f_K.

† One could split the $\rho - \omega$ mass degeneracy by adding a term of the form [10] $(TrF_{\mu\nu})^2$ to Eq. (15). Since the $\rho - \omega$ mass splitting is small (~ 2 %) , we will not consider this complication.

2.5. Total Lagrangian

Adding the kinetic energy terms for the fermions and mesons,

$$\mathcal{L}_{kin} = i\mathrm{Tr}\overline{B}\gamma_\mu\partial^\mu B + \mathrm{Tr}(\partial_\mu M^\dagger \partial^\mu M) + \frac{1}{2}\partial_\mu\chi\partial^\mu\chi - \frac{1}{4}\mathrm{Tr}(F_{\mu\nu}F^{\mu\nu}) - \frac{1}{4}\mathrm{Tr}(G_{\mu\nu}G^{\mu\nu}) \quad (18)$$

the general Lagrangian is the sum of the various terms discussed:

$$\mathcal{L} = \mathcal{L}_{kin} + \mathcal{L}_{BM} + \mathcal{L}_{BV} + \mathcal{L}_{vec} - \mathcal{V}_0 - \mathcal{V}_{SB} \qquad . \quad (19)$$

2.6. Mean field Lagrangian

To investigate the phase structure of nuclear matter at finite density we adopt the mean-field approximation (see, e.g., [7]). In this approximation scheme, the fluctuations around constant vacuum expectation values of the field operators are neglected. The fermions are treated as quantum mechanical one-particle operators. The derivative terms can be neglected and only the time-like component of the vector mesons $\omega \equiv \langle\omega_0\rangle$ and $\phi \equiv \langle\phi_0\rangle$ survive as we assume homogeneous and isotropic infinite nuclear matter. Additionally, due to parity conservation we have $\langle\pi_i\rangle = 0$. After performing these approximations, the Lagrangian (19) becomes

$$\begin{aligned}
\mathcal{L}_{BM} + \mathcal{L}_{BV} &= -\sum_i \overline{\psi_i}[g_{i\omega}\gamma_0\omega^0 + g_{i\rho}\gamma_0\tau_3\rho^0 + g_{i\phi}\gamma_0\phi^0 + m_i^*]\psi_i \\
\mathcal{L}_{vec} &= \frac{1}{2}m_\omega^2\frac{\chi^2}{\chi_0^2}\omega^2 + \frac{1}{2}m_\rho^2\frac{\chi^2}{\chi_0^2}\rho^2 + \frac{1}{2}m_\phi^2\frac{\chi^2}{\chi_0^2}\phi^2 \\
\mathcal{V}_0 &= \frac{1}{2}k_0\chi^2(\sigma^2+\zeta^2) - k_1(\sigma^2+\zeta^2)^2 - k_2(\frac{\sigma^4}{2}+\zeta^4) - k_3\chi\sigma^2\zeta \\
&\quad + k_4\chi^4 + \frac{1}{4}\chi^4\ln\frac{\chi^4}{\chi_0^4} - \frac{\delta}{3}\ln\frac{\sigma^2\zeta}{\sigma_0^2\zeta_0} \\
\mathcal{V}_{SB} &= (\frac{\chi}{\chi_0})^2[m_\pi^2 f_\pi\sigma + (\sqrt{2}m_K^2 f_K - \frac{1}{\sqrt{2}}m_\pi^2 f_\pi)\zeta] \quad ,
\end{aligned}$$

with the effective baryon masses m_i^*, which are given in Eqs. 6.

2.6.1. Grand canonical ensemble The values for the fields at a given chemical potential μ are determined by minimizing the thermodynamical potential of the grand canonical, which reads:

$$\frac{\Omega}{V} = -\mathcal{L}_{vec} + \mathcal{V}_0 + \mathcal{V}_{SB} - \mathcal{V}_{vac} - \sum_i \frac{\gamma_i}{(2\pi)^3}\int d^3k[E_i^*(k) - \mu_i^*] \quad (20)$$

The vacuum energy $\mathcal{V}_{vac}$ (the potential at $\rho = 0$) has to be subtracted in order to get a vanishing vacuum energy. γ_i denote the fermionic spin-isospin degeneracy factors ($\gamma_N = 4$, $\gamma_\Sigma = 6$, $\gamma_\Lambda = 2$, $\gamma_\Xi = 4$). The single particle energies are $E_i^*(k) = \sqrt{k_i^2 + m_i^{*2}}$

and the effective chemical potentials read $\mu_i^* = \mu_i - g_{\omega i}\omega - g_{\phi i}\phi$. The baryon chemical potentials are realated to each other by means of the nonstrange chemical potential μ_q and the strange chemical potential μ_s according to the additive quark model:

$$\mu_N = 3\mu_q \qquad \mu_\Xi = \mu_q + 2\mu_s \tag{21}$$
$$\mu_\Lambda = 2\mu_q + \mu_s = \mu_\Sigma$$

The energy density and the pressure follow from the Gibbs–Duhem relation, $\epsilon = \Omega/V + \mu_i\rho^i$ and $p = -\Omega/V$.

3. Results

3.1. Fits to nuclear matter and the hadronic spectrum

A salient feature of all chiral models are the strong vacuum constraints. In the present case they fix k_0, k_2 and k_4, in order to minimize the thermodynamical potential Ω in vacuum for given values of the fields σ_0, ζ_0 and χ_0. The parameter k_3 is fixed to the η-mass m_η. There is some freedom to vary parameters, mainly due to the unknown mass of the σ-meson, m_σ, which is determined by k_1, and due to the uncertainty of the value for the kaon decay constant f_K.
It should be noted that a reasonable nuclear matter fit with low compressibility can be found (table 1), where $m_\sigma \approx 500$ MeV (The other scalar mesons have a mass of about 1 GeV [6]). This, in the present approach, allows for an interpretation of the σ-field as the chiral partner of the π-field and as the mediator of the mid-range attractive force between nucleons, though we believe the phenomenon is in reality generated through correlated two-pion exchange [7].

$m_\pi(139)$	$m_K(495)$	$m_\eta(547)$	$m_{\eta'}(958)$
139.0	498.0	520.0	999.4
$\mathrm{m}_\rho(770)$	$\mathrm{m}_{K^*}(892)$	$\mathrm{m}_\omega(783)$	$\mathrm{m}_\phi(1020)$
783.0	857.7	783.0	1019.5
$\mathrm{m}_N(939)$	$\mathrm{m}_\Lambda(1115)$	$m_\Sigma(1193)$	$\mathrm{m}_\Xi(1315)$
939.0	1117.8	1193.1	1334.5
K [MeV]	$\frac{m_N^*}{m_N}$	f_π [MeV]	f_K [MeV]
278.7	0.62	93.0	122.0

Table 1. Fit to the hadronic mass spectrum and to nuclear matter. The results were obtained with the parameter k_0=2.3, k_1=1.4, k_2=-5.5, and k_3=-2.6 .

3.2. Finite Nuclei

If the fields are allowed to be spatial dependent, then derivative terms do not vanish and it is possible to look for the properties of finite nuclei. We adopted the Hartree formalism which is explained in [7] and included a chiral invariant quartic spin-1 meson self-interaction of the form $g_4^4 \,\mathrm{Tr}\,[(V_\mu + A_\mu)^4 + (V_\mu - A_\mu)^4]$ to allow for a fine tuning of the model.
The $SU(3)_L \times SU(3)_R \sigma$ model accounts also for a satisfactory description of finite nuclei. This is surprising, since most of the parameters are fixed to the vacuum masses, in contrast to the Walecka-model, where all the parameters are adjusted to the properties of nuclear matter or nuclei. Although no fit to nuclei is done (two parameters are fixed to nuclear matter), the agreement of the observables resulting from the model as compared to the experimental values is acceptable (table 2). Although there is no freedom to adjust the coupling constant of nucleons to the ρ meson, $g_{N\rho}$, the asymmetry energy a_4 has the reasonable value $a_4 = 39.3$ MeV.

	^{16}O		^{40}Ca		^{208}Pb	
E/A	7.53	7.98	8.15	8.55	7.65	7.86
r_{ch}	2.62	2.73	3.40	3.48	5.48	5.50
R	2.75	2.78	3.81	3.85	6.81	6.81
σ	0.81	0.84	0.93	0.98	0.89	0.90

Table 2. Bulk properties of nuclei:Prediction (left) and experimental values (right) for binding energy E/A, charge radius r_{ch}, surface tension σ and diffraction radius R of Oxigen (^{16}O), Calcium (^{40}Ca) and Lead (^{208}Pb).

3.3. Condensates and hadron masses in medium

It is instructive to see how the condensates change in the hot and dense medium, since these determine how the hadron masses change. The nonstrange condensae $\sigma \sim \langle qq \rangle$ and the strange condensate† $\zeta \sim \langle ss \rangle$) are displayed in Fig. 1 as a function of μ_q and μ_s for given temperature $T = 100$ MeV. The field σ decreases rapidly in the μ_q direction, whereas $\zeta \sim \langle ss \rangle$ doesn't change significantly at this temperature. This is different for $T = 200$ MeV. There, the nonstrange condensate‡ has already dropped to 40% of its VEV and shows only a small μ_q and μ_s dependence. This is in contrast to the

† For simplicity, we fix the gluon condensate at its vacuum expectation value ($\chi = \chi_0$).

‡ The calculation is done in the limit of vanishing coupling of the strange condensate ζ to the nucleons, i.e, for $\alpha = 1$ and $g_1^S = \sqrt{6} g_8^S$ (see Eq. 6)

strange condensate that decreased only to 80% of its VEV and reacts rather strongly for increasing μ_s .
The linear dependence of the baryon masses to the condensates causes a dropping of the effective masses (Fig. 2) with increasing temperature. This may lead to an enhanced baryon/antibaryon production. For higher chemical potentials this decrease is softened since it sets in earlier but tends towards the same high temperature limit, as can be seen comparing the curves for $\mu_q = \mu_s = 0$ and $\mu_q = 200$ MeV, $\mu_s = 0$. The 'wiggly' behaviour of the hyperon masses is due to the decrease of the nonstrange condensate at lower temperatures than the $\langle ss \rangle$.
We are currently working on the formulation of a chiral transport theory within the framework of our model which takes into account the effects of dynamical hadron masses in a heavy ion collision. These efforts may help to elucidate possible signatures of a chiral phase transition.

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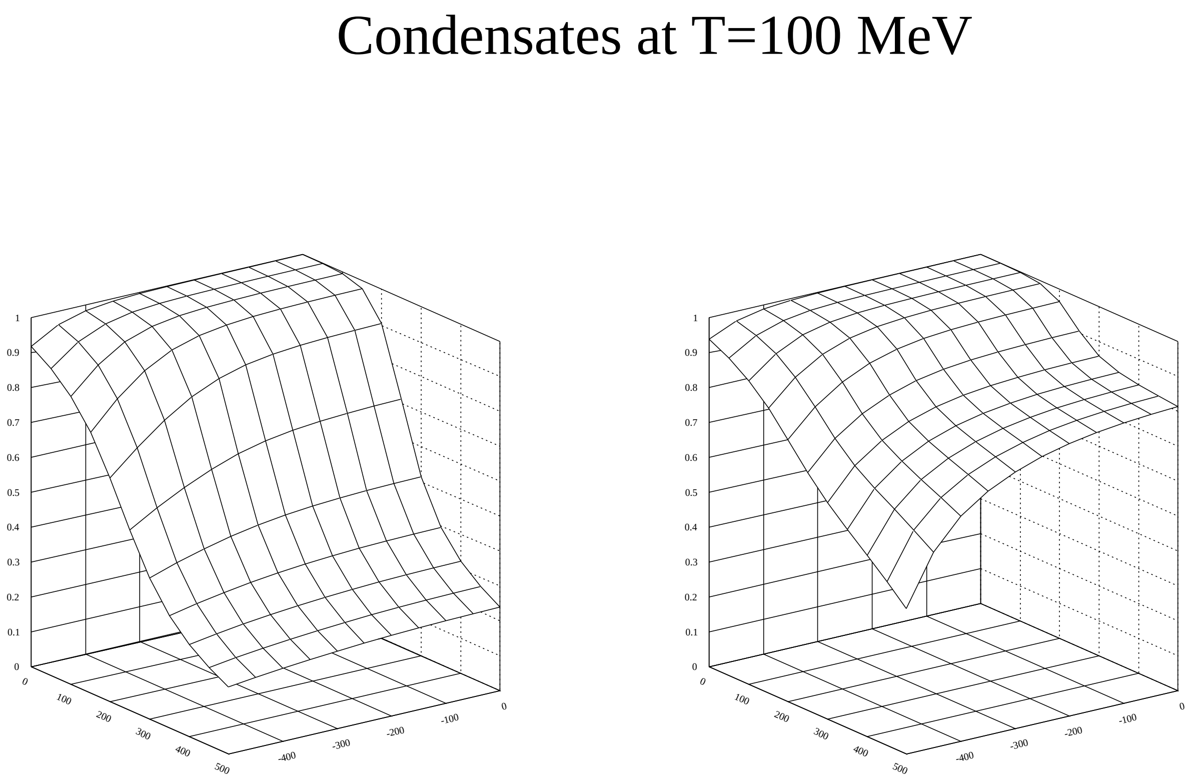

Figure 1. Nonstrange and strange condensates versus μ_q and μ_s at temperature $T = 100$ MeV.

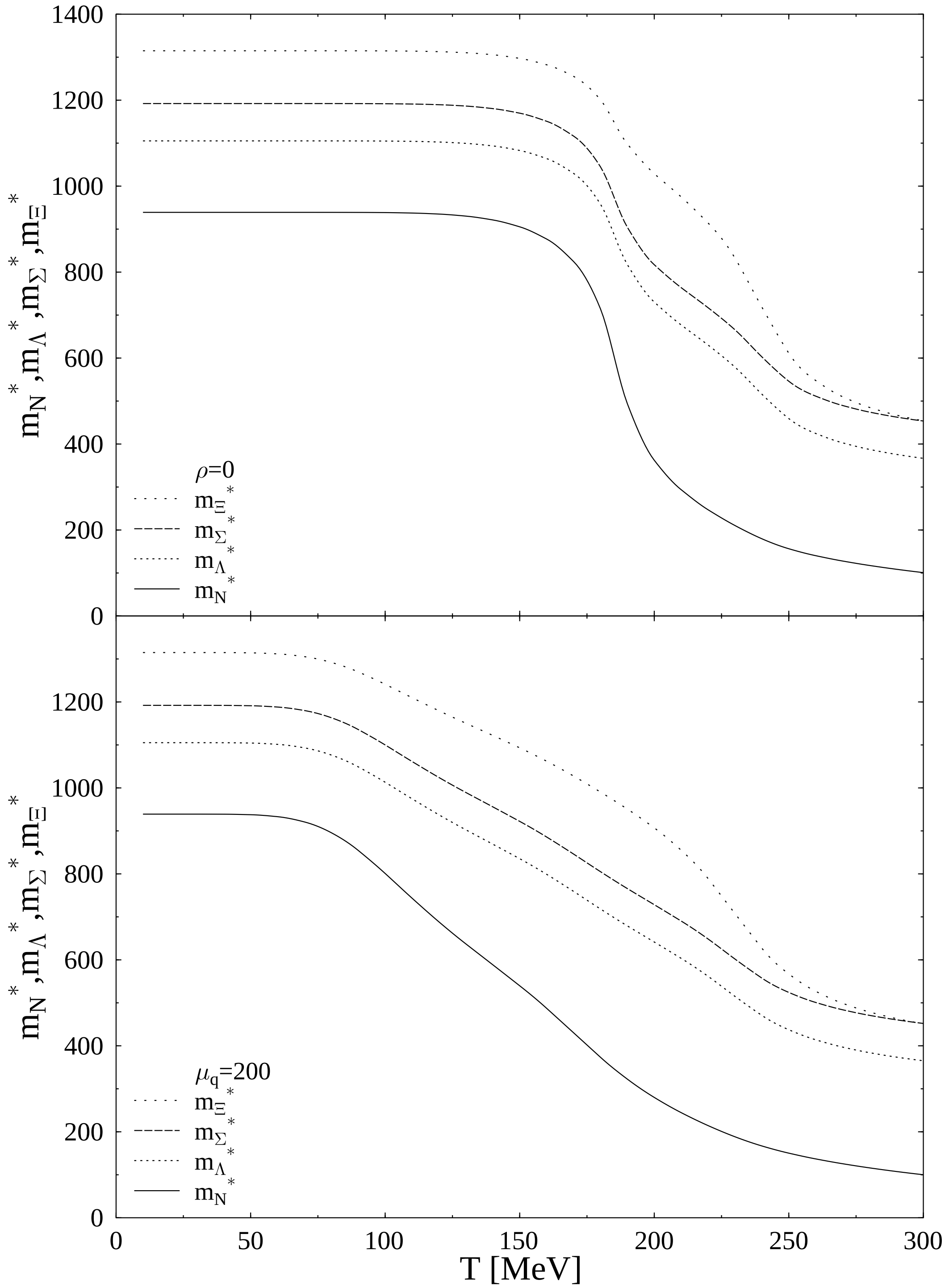

Figure 2. Baryon masses as a function of temperature for zero (above) and finite (below) chemical potentials.