The isentropic equation of state of 2-flavor QCD

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Using Taylor expansions of the pressure obtained previously in studies of 2-flavor QCD at non-zero chemical potential we calculate expansion coefficients for the energy and entropy densities up to $O(\mu_q^6)$ in the quark chemical potential. We use these series in $\mu_q/T$ to determine lines of constant entropy per baryon number ($S/N_B$) that characterize the expansion of dense matter created in heavy ion collisions. In the high temperature regime these lines are found to be well approximated by lines of constant $\mu_q/T$. In the low temperature phase, however, the quark chemical potential is found to increase with decreasing temperature. This is in accordance with resonance gas model calculations. Along the lines of constant $S/N_B$ we calculate the energy density and pressure. Within the accuracy of our present analysis we find that the ratio $p/\epsilon$ for $T > T_0$ as well as the softest point of the equation of state, $(p/\epsilon)_{\text{min}} \simeq 0.075$, show no significant dependence on $S/N_B$.

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I. INTRODUCTION

Recently studies of QCD thermodynamics on the lattice have successfully been extended to non-vanishing quark (or baryon) chemical potential using Taylor expansions\cite{1,2} as well as reweighting techniques\cite{3} around the limit of vanishing quark chemical potential ($\mu_q = 0$) and analytic continuations of numerical calculations performed with an imaginary chemical potential\cite{4}. The complementary approach at fixed baryon number\cite{5,6} has also been used recently in simulations of 4-flavor QCD with light quarks\cite{7}. These calculations yield the additional contribution to the pressure that arises from the presence of a net excess of baryons over anti-baryons in a strongly interacting medium. We will further explore here the approach based on a Taylor expansion of the pressure and make use of our earlier calculations for 2-flavor QCD\cite{2} to determine also the energy and entropy densities at non-vanishing baryon chemical potential.

In a heavy ion collision a dense medium is created which after thermalization is expected to expand without further generation of entropy ($S$) and with fixed baryon number ($N_B$) or, equivalently, with fixed quark number $N_q = 3N_B$. During the isentropic expansion the ratio $S/N_B$ thus remains constant\footnote{Here $N_B$ denotes the net baryon number, i.e. the excess of baryons over anti-baryons.}. The cooling of the expanding system then is controlled by the equation of state on lines of constant $S/N_B$. From a knowledge of the energy density and pressure at non-vanishing quark chemical potential we can calculate trajectories in the $\mu_q$-$T$ phase diagram of QCD that correspond to constant $S/N_B$ and can determine the isentropic equation of state on these trajectories.

This paper is organized as follows. In the next section we summarize the basic setup for calculating Taylor expansions for bulk thermodynamic observables and present a calculation of the Taylor expansion coefficients for energy and entropy densities at non-vanishing quark chemical potential up to $O(\mu_q^6)$. We use these results in Section III to determine lines of constant $S/N_B$ and study the temperature dependence of pressure and energy density along these lines. Our conclusions are given in section IV.

II. TAYLOR EXPANSION OF PRESSURE, ENERGY AND ENTROPY DENSITY

The analysis we are going to present in this paper is based on numerical results previously obtained in simulations of 2-flavor QCD on lattices of size $16^3 \times 4$ with an improved staggered fermion action\cite{1,2}. In these calculations Taylor expansion coefficients for the pressure have been obtained up to $O(\mu_q^6)$ for a fixed bare quark mass value $(\hat{m} = 0.1)$ which at temperatures close to the transition temperature $(T_0)$ corresponds to a still quite large pion mass of about...
770 MeV. However, in particular at temperatures above the QCD transition temperature the remaining quark mass dependence of thermodynamic observables is nonetheless small as deviations from the massless limit are controlled by the quark mass in units of the temperature, which is a small number.

We closely follow here the approach and notation used in Ref. \[2\]. We start with a Taylor expansion for the pressure in 2-flavor QCD for non-vanishing quark chemical potential \(\mu_q\) (and vanishing isospin chemical potential, \(\mu_I \equiv 0\)),

\[
\frac{p}{T^4} = \frac{1}{V T^3} \ln Z(T, \mu_q) = \sum_{n=0}^{\infty} c_n(T, m_q) \left(\frac{\mu_q}{T}\right)^n ,
\]

(1)

with

\[
c_n(T, m_q) = \frac{1}{n!} \frac{\partial^n \ln Z(T, \mu_q)}{\partial (\mu_q/T)^n} \bigg|_{\mu_q=0} .
\]

(2)

Here we explicitly indicated that the Taylor expansion coefficients \(c_n\) depend on temperature as well as the quark mass\(^2\). In the context of lattice calculations it is more customary to think of this quark mass dependence in terms of a dependence on an appropriately chosen ratio of hadron masses characterizing lines of constant physics. We thus may replace \(m_q\) by a ratio of pseudo-scalar (pion) and vector (rho) meson masses, \(m_q \equiv (m_{PS}/m_V)^2\).

Due to the invariance of the partition function under exchange of particle and anti-particle the Taylor expansion is a series in even powers of \(\mu_q/T\); expansion coefficients \(c_n\) vanish for odd values of \(n\). From the pressure we immediately obtain the quark number density,

\[
\frac{n_q}{T^3} = \frac{1}{V T^3} \frac{\partial \ln Z(T, \mu_q)}{\partial \mu_q/T} \bigg|_T = \sum_{n=2}^{\infty} n c_n(T, m_q) \left(\frac{\mu_q}{T}\right)^{n-1} .
\]

(3)

Using standard thermodynamic relations we also can calculate the difference between energy density \((\epsilon)\) and three times the pressure,

\[
\frac{\epsilon - 3p}{T^4} = \sum_{n=0}^{\infty} c_n'(T, m_q) \left(\frac{\mu_q}{T}\right)^n ,
\]

(4)

where

\[
c_n'(T, m_q) = T \frac{d c_n(T, m_q)}{dT} .
\]

(5)

Combining Eqs. (1) and (3) we then obtain Taylor expansions for the energy and entropy densities,

\[
\frac{\epsilon}{T^4} = \sum_{n=0}^{\infty} (3c_n(T, m_q) + c_n'(T, m_q)) \left(\frac{\mu_q}{T}\right)^n ,
\]

\[
\frac{s}{T^3} = \frac{\epsilon + p - \mu_q n_q}{T^4} = \sum_{n=0}^{\infty} ((4-n)c_n(T, m_q) + c_n'(T, m_q)) \left(\frac{\mu_q}{T}\right)^n .
\]

(6)

The expansion coefficients \(c_n(T, m_q)\) have been calculated in 2-flavor QCD at several values of the temperature and for a fixed value of the bare quark mass, \(\hat{m} = 0.1\). We note that the bare coupling \(\hat{m}\) introduces an implicit temperature dependence if \(\hat{m}\) is kept fixed while the lattice cut-off is varied. The latter controls the temperature of the system on lattices of temporal extent \(N_t\), i.e. \(T^{-1} = N_t a\). The difference between derivatives taken at fixed \(\hat{m}\) and fixed \(m_q\), however, is negligible at high temperature, where the quark mass dependence of bulk thermodynamic quantities is \(O(m_q/T)^2\) and, of course, vanishes at all temperatures in the chiral limit as this also defines a line of constant physics. We thus approximate the derivative, Eq. (5) by a derivative taken at fixed \(\hat{m}\) which becomes exact in the chiral limit. These derivatives are evaluated using finite difference approximants. Alternatively we may express

\[\text{---}\]

\(^2\) We assume that the thermodynamic limit has been taken so that we can ignore any volume dependence of thermodynamic quantities. For an aspect ratio \(TV^{1/3} = 4\) and the quark mass value used in our calculation finite volume effects are expected to be small.
derivatives of Taylor expansion coefficients with respect to $T$ in terms of derivatives with respect to the bare lattice couplings $\beta = 6/g^2$ and $\hat{m}$ and rewrite Eq. 5 as,

$$c'_n(T, m_q) = -\frac{a \partial c_n(T, m_q)}{a \partial \beta} - \frac{a \partial m \partial c_n(T, m_q)}{a \partial m}.$$ \hfill (7)

Here the two lattice $\beta$-functions, $ad\beta/da$ and $ad\hat{m}/da$, have to be determined by demanding that the temperature derivative is taken along lines of constant physics \cite{10, 11}. The $\beta$-function controlling the variation of the bare quark mass with the lattice cut-off, $a \partial \hat{m}/da$, is proportional to the bare quark mass $\hat{m}$ and thus vanishes in the chiral limit \cite{10}. For small quark masses the first term in Eq. 4 thus gives the dominant contribution to $c'_n(T, m_q)$. We have evaluated this first term in Eq. 4 for $n = 2$ and find that this approximation for $c'_2$ agrees within errors with the calculation of $c'_2$ from finite difference approximants. Only close to $T_0$ we find differences of the order of 10%.

The expansion coefficients $c_n(T, m_q)$ have been calculated in \cite{2} for $\hat{m} = 0.1$ at a set of 16 temperature values in the interval $T/T_0 \in [0.81, 1.98]$. From this set of expansion coefficients we have calculated the partial derivatives $c'_n$ at temperature $T$ by averaging over finite difference approximants for left and right derivatives at $T$. With this we have constructed the expansion coefficients for the energy and entropy densities,

$$c_n (T, m_q) = 3c_n(T, m_q) + c'_n(T, m_q),$$

$$s_n (T, m_q) = (4 - n)c_n(T, m_q) + c'_n(T, m_q).$$ \hfill (8)

Results for the 2nd, 4th and 6th order expansion coefficients of energy and entropy densities obtained in the approximation discussed above are shown in Fig. 11. Also shown in this figure are the corresponding expansion coefficients for the pressure ($p_n \equiv c_n$). For temperatures larger than $T/T_0 \simeq 1.5$ all expansion coefficients satisfy quite well the ideal gas relations, $\epsilon^{SB}_n = 3p_n^{SB}$ and $c^{SB}_n = 3s^{SB}_n/2$. As noted already in Ref. \cite{2} results for the 2nd and 4th order expansion coefficients are close to those of an ideal Fermi gas which describes the high temperature limit for the energy density of 2-flavor QCD,

$$\epsilon^{SB}_n = \frac{7\pi^2}{10} + 3 \left( \frac{\mu_q}{T} \right)^2 + \frac{3}{2\pi^2} \left( \frac{\mu_q}{T} \right)^4.$$ \hfill (9)

Nonetheless, quantitative statements on the approach to ideal gas behavior may be at present premature. Deviations from the continuum result clearly arise in our current analysis which has been performed on lattices with a lattice spacing in units of the inverse temperature given by $aT \equiv 1/N_\tau = 1/4$. Although cut-off effects are reduced due the usage of an improved action the approach to the continuum limit has to be analyzed further through calculations on lattices with larger temporal extent $N_\tau$.

Although errors on the 6th order expansion coefficients are large it is apparent that subsequent orders in the expansion increase in magnitude thus making the expansions for energy and entropy density well behaved for $\mu_q/T \lesssim 1$. In fact, the 6th order contribution to $\epsilon' T^4$ and $s/T^3$ never exceeds 10% for $\mu_q/T < 0.9$. We will see below that this is the regime of interest for a discussion of the thermodynamics of matter created in heavy ion collisions at the SPS and RHIC and even is appropriate for the energy range to be covered by the AGS and the future FAIR facility at GSI/Darmstadt. Results for the $\mu_q$-dependent contribution to the energy and entropy densities, $\Delta \epsilon'/T^4 \equiv \epsilon'/T^4 - \epsilon_0$ and $\Delta s/T^3 \equiv s/T^3 - s_0$, obtained in 4th order Taylor expansion are shown in Fig. 2 for some values of $\mu_q/T$.

### III. LINES OF CONSTANT ENTROPY PER BARYON NUMBER

When discussing the equation of state at finite baryon number density we have to eliminate the quark chemical potential in the thermodynamic quantities introduced in the previous section. We may do this by demanding that a thermodynamic quantity, for instance the quark number density, stays constant. We will use here a prescription more appropriate for discussing the thermodynamics of matter created in relativistic heavy ion collisions. After equilibration the dense medium created in such a collision will expand along lines of constant entropy per baryon. It then is of interest to calculate thermodynamic quantities along such isentropic lines.

We thus should first analyze how the quark chemical potential changes along lines of constant $S/N_B$. It is readily seen that in an ideal quark-gluon gas these lines are defined by $\mu_q/T = \text{const}$,

$$S = 3 \frac{37\pi^2}{45} + \left( \frac{\mu_q}{T} \right)^2.$$ \hfill (10)

In the zero temperature limit, however, the resonance gas reduces to a degenerate Fermi gas of nucleons and the chemical potential approaches a finite value to obtain finite baryon number and entropy, i.e. $\mu_q/T \sim 1/T$. 

FIG. 1: The 2nd, 4th and 6th order Taylor expansion coefficients for pressure ($p_n \equiv c_n$), energy density ($\epsilon_n$) and entropy density ($s_n$) as functions of $T/T_0$.

In Fig. 2 we show the values for $\mu_q$ needed to keep $S/N_B$ fixed. These lines of constant $S/N_B$ in the $T$-$\mu_q$ plane have been obtained by calculating the total entropy density for 2-flavor QCD from Eq. 6 using results for the pressure and energy density calculated at $\mu_q = 0$ [9, 12] and the corresponding $\mu_q$ dependent contributions shown in Fig. 1. The ratio of $s/T^3$ and $n_q/T^3$ obtained in this way is then solved numerically for $\mu_q/T$.

We find that isentropic expansion at high temperature indeed is well represented by lines of constant $\mu_q/T$ down to temperatures close to the transition, $T \simeq 1.2T_0$. In the low temperature regime we observe a bending of the isentropic lines in accordance with the expected asymptotic low temperature behavior. The isentropic expansion lines for matter created at SPS correspond to $S/N_B \simeq 45$ while the isentropes at RHIC correspond to $S/N_B \simeq 300$. The energy range
FIG. 3: Lines of constant entropy per quark number versus $\mu_q/T$ (left) and in physical units using $T_0 = 175$ MeV to set the scales (right). In the left hand figure we show results obtained using a $4^{th}$ (full symbols) and $6^{th}$ (open symbols) order Taylor expansion of the pressure, respectively. Data points correspond to $S/N_B = 300, 150, 90, 60, 45, 30$ (from left to right). The vertical lines indicate the corresponding ideal gas results, $\mu_q/T = 0.08, 0.16, 0.27, 0.41, 0.54$ and $0.82$ in decreasing order of values for $S/N_B$. For a detailed description of the right hand figure see the discussion given in the text.

FIG. 4: The complete energy density (left) and pressure (right) evaluated on lines of constant $S/N_B$ using Taylor expansions up to $6^{th}$ order in the quark chemical potential. Shown are results for $S/N_B = 30, 45, 300$.

of the AGS which also corresponds to an energy range relevant for future experiments at FAIR/Darmstadt is well described by $S/N_B \simeq 30$. These lines are shown in Fig. 3 (right) together with points characterizing the chemical freeze-out of hadrons measured at AGS, SPS and RHIC energies. These points have been obtained by comparing experimental results for yields of various hadron species with hadron abundances in a resonance gas [13, 14]. The solid curve shows a phenomenological parametrization of these freeze-out data [14]. In general our findings for lines of constant $S/N_B$ are in good agreement with phenomenological model calculations that are based on combinations of ideal gas and resonance gas equations of state at high and low temperature, respectively [15, 16].

Our current analysis yields stable results for lines of constant $S/N_B$ also for temperatures $T < \sim T_0$ and values of the chemical potential $\mu_B \simeq 400$ MeV. This value of $\mu_B$ is larger than recent estimates for the location of the chiral critical point in 2-flavor [17] and (2+1)-flavor QCD. We thus may expect modifications to our current analysis at temperatures below $T_0$ once we include higher orders in the Taylor expansion and/or perform this analysis at smaller values of the quark mass.

We now can proceed and calculate energy density and pressure on lines of constant entropy per baryon number using our Taylor expansion results up to $O(\mu_q^6)$. In Fig. 4 we show both quantities for the parameters relevant for AGS (FAIR), SPS and RHIC energies. At high temperatures the relevant values of the quark chemical potential in an ideal quark gas are $\mu_q/T = 0.82, 0.54$ and $0.08$, respectively. However, as can be seen from Fig. 3 at temperatures $T \lesssim 1.8 T_0$ the required values for the chemical potentials are significantly smaller. In particular, for the AGS (FAIR) energies ($S/N_B = 30$) we find $\mu_q/T = 0.77$ at $T \simeq 1.8 T_0$. The $6^{th}$ order Taylor expansion thus is still well behaved.
at these values of the quark chemical potential. The dependence of $\epsilon$ and $p$ on $S/N_B$ cancels to a large extent in the ratio $p/\epsilon$, which is most relevant for the analysis of the hydrodynamic expansion of dense matter. This may be seen by considering the leading $O(\mu_q^2)$ correction,

$$
\frac{p}{\epsilon} = \frac{1}{3} - \frac{1}{3} \frac{\epsilon_0 - 3p_0}{\epsilon_0} \left( 1 + \left[ \frac{c_2}{\epsilon_0 - 3p_0} - \frac{c_2}{\epsilon_0} \right] \left( \frac{\mu_q}{T} \right)^2 \right).
$$

(11)

In Fig. 5 we show $p/\epsilon$ as function of $T/T_0$ (left) as well as function of the energy density (right). In the latter case we again used $T_0 = 175$ MeV to set the scale. The small dependence on $S/N_B$ visible in Fig. 5(left) gets further reduced when considering the ratio $p/\epsilon$ at fixed energy density (Fig. 5(right)). The softest point of the equation of state is found to be $(p/\epsilon)_{\text{min}} \approx 0.075$ and within our current numerical accuracy it is independent of $S/N_B$. However, the analysis of the temperature regime $T \lesssim T_0$ at present clearly suffers from poor statistics and still needs a more detailed analysis. We also find that the equation of state for $T_0 < T < 2T_0$ is well parametrized by

$$
\frac{p}{\epsilon} = \frac{1}{3} \left( 1 - \frac{1.2}{1 + 0.5\epsilon \text{ fm}^3/\text{GeV}} \right).
$$

(12)

We note, however, that this phenomenological parametrization is not correct at asymptotically large temperatures as it is obvious that Eq. (12) would lead to corrections to the ideal gas result $p/\epsilon = 1/3$ that are proportional to $T^{-4}$ while perturbative corrections vanish only logarithmically as function of $T$.

IV. CONCLUSIONS

We have determined Taylor expansion coefficients for the energy and entropy densities in 2-flavor QCD at non-zero quark chemical potential. At present these expansion coefficients have been determined from calculations performed at one value of the bare quark mass which gives a good description of QCD thermodynamics at high temperature. In this temperature regime lines of constant entropy per baryon are well described by lines of constant $\mu_q/T$. On these isentropic lines we have determined the equation of state and find that the ratio of pressure and energy density shows remarkably little dependence on the ratio $S/N_B$.

The regime close to and below the transition temperature $T_0$ is not yet well controlled in our current analysis. Smaller quark masses and the introduction of a non-vanishing strange quark mass, kept fixed in physical units, will be needed to explore this regime as well as the approach to the 3-flavor high temperature limit in more detail.

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[1] C. R. Allton, S. Ejiri, S. J. Hands, O. Kaczmarek, F. Karsch, E. Laermann and C. Schmidt, Phys. Rev. D 68 (2003) 014507.
[2] C. R. Allton, M. Döring, S. Ejiri, S. J. Hands, O. Kaczmarek, F. Karsch, E. Laermann, K. Redlich, Phys. Rev D 71 (2005) 054508.
[3] Z. Fodor, S. D. Katz and K. K. Szabo, Phys. Lett. B 568 (2003) 73;
F. Csikor, G. I. Egri, Z. Fodor, S. D. Katz, K. K. Szabo and A. I. Tóth, JHEP 0405 (2004) 046.
[4] M. D’Elia and M. P. Lombardo, Phys. Rev. D 70 (2004) 074509.
[5] D.E. Miller and K. Redlich, Phys. Rev. D 35 (1987) 2524.
[6] J. Engels, O. Kaczmarek, F. Karsch and E. Laermann, Nucl. Phys. B 558 (1999) 307.
[7] S. Kratochvila and Ph. de Forcrand, PoS LAT2005 (2005) 167.
[8] C. R. Allton, S. Ejiri, S. J. Hands, O. Kaczmarek, F. Karsch, E. Laermann, Ch. Schmidt and L. Scorzato, Phys. Rev. D 66 (2002) 074507.
[9] F. Karsch, E. Laermann, A. Peikert, Nucl. Phys. B 605 (2001) 579.
[10] T. Blum, L. Kärkäinen, D. Toussaint and S. Gottlieb, Phys. Rev. D 51 (1995) 5153.
[11] A. AliKhan et al. (CP-PACS), Phys. Rev. D 64 (2001) 074510.
[12] A. Peikert, PhD thesis, Bielefeld, May 2000.
[13] J. Cleymans and K. Redlich, Phys. Rev. Lett. 81 (1998) 5284
J. Cleymans and K. Redlich, Phys. Rev. C 60 (1999) 054908.
[14] J. Cleymans, H. Oeschler and K. Redlich, S. Wheaton, hep-ph/0511094
[15] C. M. Hung and E. Shuryak, Phys. Rev. C 57 (1998) 1891.
[16] V. D. Toneev, J. Cleymans, E. G. Nikonov, K. Redlich and A. A. Shanenko, J. Phys. G 27 (2001) 827, nucl-th/0011029
[17] R.V. Gavai and S. Gupta, Phys. Rev. D 71 (2005) 114014.
[18] Z. Fodor and S. D. Katz, JHEP 0404 (2004) 050.