Evidence for In-Medium Changes of Four-Quark Condensates

R. Thomas,[1] S. Zschocke,[1,2] and B. Kämpfer[1,3]

[1] Institut für Kern- und Hadronenphysik, Forschungszentrum Rossendorf, PF 510119, 01314 Dresden, Germany
[2] Theoretical and Computational Physics Section, University of Bergen, 5007 Bergen, Norway
[3] Institut für Theoretische Physik, Technische Universität Dresden, 01062 Dresden, Germany

(Received 11 July 2005; published 28 November 2005)

Utilizing the QCD sum rule approach to the behavior of the ω meson in nuclear matter we derive evidence for in-medium changes of particular four-quark condensates from the recent CB-TAPS experiment for the reaction $\gamma + \mathrm{A} \rightarrow \mathrm{A}' + \omega(\rightarrow \pi^0\gamma)$ with $A = Nb$ and LH_2.

DOI: 10.1103/PhysRevLett.95.232301 PACS numbers: 24.85.+p, 12.38.Lg, 12.40.Yx

The chiral condensate $\langle\bar{q}q\rangle$ is an order parameter for the spontaneous breaking of chiral symmetry in the theory of strong interaction (cf. e.g., Ref. [1] for introducing this topic). The role of $\langle\bar{q}q\rangle$ is highlighted, e.g., by the Gell-Mann–Oakes–Renner relation $m_\pi^2 f_\pi^2 \propto -\langle\bar{q}q\rangle$ (cf. Ref. [2]; the explicit chiral symmetry breaking is essential for a finite pion mass m_π, while the relation of the pion decay constant f_π to $\langle\bar{q}q\rangle$ qualifies the latter as an order parameter) or by Ioffe's formula $M_N \propto -\langle\bar{q}q\rangle$ for the nucleon mass (cf. Ref. [3] and in particular the discussion in Ref. [4]). There is growing evidence that the quark-gluon condensate is another order parameter [5]. The QCD trace anomaly related to scale invariance breaking gives rise to the gluon condensate. There are many other condensates characterizing the complicated structure of the QCD vacuum. In a medium, described by temperature and baryon density n, these condensates change; i.e., the ground state is rearranged. Since hadrons are considered as excitations above the vacuum, a vacuum change should manifest itself as a change of the hadronic excitation spectrum. This idea triggered widespread activities to search for in-medium modifications of hadrons. Such in-medium modifications of hadronic observables are found (cf. the lists in Ref. [6,7]), and it is timely to relate them to corresponding order parameters.

We deduce here evidence for a noticeable drop of in-medium four-quark condensates in cold nuclear matter from results of the recent CB-TAPS experiment [6] for the reaction $\gamma + A \rightarrow A' + \omega(\rightarrow \pi^0\gamma)$. The CB-TAPS collaboration observed the occurrence of additional low-energy ω decay strength for a Nb ($A = 93$) target compared to a LH_2 ($A = 1$) target. The link of observables to quark and gluon condensates is established by QCD sum rules [8], which are expected to be sensitive to four-quark condensates in the vector channels [9]. Four-quark condensate combinations which contain only left-right helicity flipping terms (as the chiral condensate does) represent other order parameters of chiral symmetry.

Concentrating on the isoscalar part of the causal current-current correlator [3]

$$\Pi^\omega(q, n) = \frac{i}{3}\int d^4x e^{iqx}\langle\Omega|\mathcal{T} j_\mu^\omega(x) j^{\omega\mu}(0)|\Omega\rangle, \qquad (1)$$

here for the ω meson with the current $j_\mu^\omega = (\bar{u}\gamma_\mu u + \bar{d}\gamma_\mu d)/2$ and nuclear matter states $|\Omega\rangle$ (the symbol $\mathcal{T}$ means time ordering, and u, d denote quark field operators), an operator product expansion and a Borel transformation (cf. Refs. [3,10] for arguments in favor of Borel sum rules) of the twice-subtracted dispersion relation result in

$$\Pi^\omega(0, n) - \frac{1}{\pi}\int_0^\infty ds \frac{\mathrm{Im}\Pi^\omega(s, n)}{s} e^{-s/\mathcal{M}^2} = c_0\mathcal{M}^2 + \sum_{j=1}^\infty \frac{c_j}{(j-1)!\mathcal{M}^{2(j-1)}}, \qquad (2)$$

where $\Pi^\omega(0, n) = 9n/(4M_N)$ with the nucleon mass M_N is a subtraction constant having the meaning of Landau damping or ωN forward scattering amplitude, and the coefficients c_j contain condensates and Wilson coefficients; $\mathcal{M}$ is the Borel mass. The first coefficients c_j have been spelled out in many papers (cf. Ref. [11] for our notation, and Ref. [12] for an anomalous contribution) and are not reproduced here in full length. $c_0 = (1 + \frac{\alpha_s}{\pi})/(8\pi^2)$ is the perturbative term. $c_1 \propto m_q^2$ is exceedingly small due to the small current quark mass m_q. In c_2 the gluon condensate (being less sensitive to medium effects), some moments of the parton distribution in the nucleon (combined with a density dependence), and the renormalization group invariant combination $m_q\langle\bar{q}q\rangle$ (being numerically tiny) enter. The latter fact makes the Borel sum rule insensitive to the genuine chiral condensate, but sensitive to four-quark condensates which enter c_3, among other quantities related to expectation values of certain traceless and symmetric twist-2 and twist-4 operators. To be specific, the flavor-mixing condensates $\frac{2}{9}\langle\bar{u}\gamma^\mu\lambda_A u\bar{d}\gamma_\mu\lambda_A d\rangle + \langle\bar{u}\gamma_5\gamma^\mu\lambda_A u\bar{d}\gamma_5\gamma_\mu\lambda_A d\rangle$ and the pure flavor four-quark condensates (for which we employ $u - d$ isospin symmetry; γ_μ and λ_A stand for Dirac and Gell-Mann matrices) $\frac{2}{9}\langle\bar{q}\gamma^\mu\lambda_A q\bar{q}\gamma_\mu\lambda_A q\rangle + \langle\bar{q}\gamma_5\gamma^\mu\lambda_A q\bar{q}\gamma_5\gamma_\mu\lambda_A q\rangle$ enter c_3. (c_4 will be discussed below.) Our strategy to deal with these condensates is as follows: (i) the factorized expressions (which might fail badly [3,12]) are corrected by factors

κ_Ω (with Ω being a label of the respective four-quark condensate) using $\langle \bar{u}\gamma^\mu \lambda_A u \bar{d}\gamma_\mu \lambda_A d\rangle = -\kappa_1 \frac{4}{9\pi^2}\frac{Q_0^2}{f_\pi^2}\langle\bar{q}q\rangle^2$, $\langle \bar{u}\gamma_5\gamma^\mu \lambda_A u \bar{d}\gamma_5\gamma_\mu \lambda_A d\rangle = \kappa_2 \frac{4}{9\pi^2}\frac{Q_0^2}{f_\pi^2}\langle\bar{q}q\rangle^2$ (where Q_0 is a cutoff related to the $\rho - \omega$ mass splitting; both expressions are already beyond the ground state saturation [11]), $\langle \bar{q}\gamma^\mu \lambda_A q \bar{q}\gamma_\mu \lambda_A q\rangle = -\frac{16}{9}\kappa_3\langle\bar{q}q\rangle^2$, $\langle \bar{q}\gamma_5\gamma^\mu \lambda_A q \bar{q}\gamma_5\gamma_\mu \lambda_A q\rangle = \frac{16}{9}\kappa_4\langle\bar{q}q\rangle^2$; $\kappa_{1,2} = 0$ and $\kappa_{3,4} = 1$ recover the factorized terms in the ground state saturation approximation; (ii) expand κ_Ω in density [13], i.e., $\kappa_\Omega = \kappa_\Omega^{(0)} + \kappa_\Omega^{(1)} n$, use the known sigma term σ_N in $\langle\bar{q}q\rangle = \langle\bar{q}q\rangle_0 + \xi n$ with $\xi = \sigma_N/(2m_q)$ [3], and linearize the resulting expressions [14]; (iii) add up all contributions with their corresponding prefactors to get a common factor $\kappa_0 = \frac{9}{28\pi^2}\frac{-Q_0^2}{f_\pi^2}(\frac{2}{9}\kappa_1^{(0)} - \kappa_2^{(0)}) - \frac{2}{7}\kappa_3^{(0)} + \frac{9}{7}\kappa_4^{(0)}$ for the vacuum contribution, $-\frac{112}{81}\pi\alpha_s\kappa_0\langle\bar{q}q\rangle_0^2$, and a common factor $\kappa_N = \kappa_0 + \frac{\langle\bar{q}q\rangle_0}{2\xi}(\frac{9}{28\pi^2}\frac{-Q_0^2}{f_\pi^2}(\frac{2}{9}\kappa_1^{(1)} - \kappa_2^{(1)}) - \frac{2}{7}\kappa_3^{(1)} + \frac{9}{7}\kappa_4^{(1)})$ for the density dependent medium contribution of the mentioned four-quark condensates [cf. the first term in Eq. (10) below]. κ_0 enters the vacuum sum rule and has to be adjusted properly with other quantities to get the correct vacuum ω mass, while κ_N is the subject of our further consideration. Because of the mixing of density dependencies of κ_Ω and $\langle\bar{q}q\rangle$ even the accurate knowledge of κ_0 does not fix κ_N.

No density dependence of the four-quark condensates would imply $\kappa_N = 0$, while strong density dependencies will result in a sizeable value of κ_N, unless the terms contributing to κ_N cancel. The estimate in Ref. [11] points to small values of $\kappa_{1,2}^{(1)}$ thus having essentially the density dependence of the combined pure flavor scalar dimension-6 condensates $\frac{2}{9}\langle \bar{q}\gamma^\mu \lambda_A q \bar{q}\gamma_\mu \lambda_A q\rangle + \langle \bar{q}\gamma_5\gamma^\mu \lambda_A q \bar{q}\gamma_5\gamma_\mu \lambda_A q\rangle$ to be constrained.

Large-N_c arguments [3,15] favor $\kappa_N = \kappa_0$. Previously, often the factorization $\langle\bar{q}\cdots q\bar{q}\cdots q\rangle \to \langle\bar{q}q\rangle^2$ has been used. Here, we study explicitly, however, the role of the four-quark condensates using the square of the genuine chiral condensate only to set the scale, as outlined above. The integral in the left hand side of (2) can be decomposed in a low-lying resonance part, $\int_0^{s_\omega} ds\, \mathrm{Im}\Pi^\omega(s,n)s^{-1}e^{-s/\mathcal{M}^2}$, and the continuum part, $\int_{s_\omega}^\infty ds\, \mathrm{Im}\Pi^\omega(s,n)s^{-1}e^{-s/\mathcal{M}^2} \equiv -\pi\mathcal{M}^2 c_0 e^{-s_\omega/\mathcal{M}^2}$, both depending on the continuum threshold s_ω. The quantity

$$m_\omega^2(n,\mathcal{M}^2,s_\omega) \equiv \frac{\int_0^{s_\omega} ds\, \mathrm{Im}\Pi^\omega(s,n)e^{-s/\mathcal{M}^2}}{\int_0^{s_\omega} ds\, \mathrm{Im}\Pi^\omega(s,n)s^{-1}e^{-s/\mathcal{M}^2}} \tag{3}$$

is a normalized moment with s meaning the coordinate of the center of gravity of $\mathrm{Im}\Pi^\omega(s,n)e^{-s/\mathcal{M}^2}/s$ in the interval $s = 0\cdots s_\omega$. Clearly, when additional strength of $\mathrm{Im}\Pi^\omega$ at lower values of s is caused by in-medium effects as observed in Ref. [6], then the center of gravity shifts to the left; i.e., m_ω^2 becomes smaller. Direct use of the count rates in the middle panel of Fig. 2 in Ref. [6] as estimator of $\mathrm{Im}\Pi^\omega$ in the interval $s = 0.41\cdots 0.77$ GeV^2 yields $m_\omega^2(\mathrm{LH}_2) = 0.599$ GeV^2 and $m_\omega^2(Nb) = 0.568$ GeV^2 for $\mathcal{M} \sim \mathcal{O}(1)$ GeV. Instead of testing the consistency of a particular model for $\mathrm{Im}\Pi^\omega(s,n)$ with the sum rule, we suggest here to use the experimental information on $\mathrm{Im}\Pi^\omega$ to find constraints on the QCD side of the sum rule. In fact, the ω decay rate $\omega \to \pi^0\gamma$ is given by $dR_{\omega\to\pi^0\gamma}/d^4q = (6d/f_\pi)^2(\pi/[3q^2]) \times (q^2 - m_\pi^2)^3\, \mathrm{Im}\Pi^\omega(q^2 = s)$ with $d = 0.011$. However, acceptance and efficiency corrections to the results of Ref. [6] need to be invoked and the fraction of events, where the rate $dR_{\omega\to\pi^0\gamma}/dM_{\pi^0\gamma}$ is shifted to smaller values of $M_{\pi^0\gamma}$ (being the invariant mass of the π^0 and γ decay products of ω) by final state interaction of the decay π^0 in the ambient nuclear medium [16], must be corrected for as well. We postpone such a quantitative and model dependent study for future work and consider qualitatively here the implication of the observation of Ref. [6], i.e., the occurrence of additional ω decay strength at $M_{\pi^0\gamma} < m_\omega^{(0)}$ which translates into $m_\omega < m_\omega^{(0)} = 0.782$ GeV for low-momentum ω decaying in the Nb nucleus.

With (3) the truncated QCD sum rule (2) for the ω meson can be arranged as [11]

$$m_\omega^2(n,\mathcal{M}^2,s_\omega) = \frac{c_0\mathcal{M}^2[1-(1+\frac{s_\omega}{\mathcal{M}^2})e^{-s_\omega/\mathcal{M}^2}] - \frac{c_2}{\mathcal{M}^2} - \frac{c_3}{\mathcal{M}^4} - \frac{c_4}{2\mathcal{M}^6}}{c_0(1-e^{-s_\omega/\mathcal{M}^2}) + \frac{c_1}{\mathcal{M}^2} + \frac{c_2}{\mathcal{M}^4} + \frac{c_3}{2\mathcal{M}^6} + \frac{c_4}{6\mathcal{M}^8} - \frac{\Pi^\omega(0,n)}{\mathcal{M}^2}}. \tag{4}$$

This sum rule is to be handled as usual (cf. Refs. [10,11]): determine the sliding Borel window by requiring that (i) the sum of the $c_{3,4}$ terms in Eq. (2) does not contribute more than 10% to the right hand side; (ii) the continuum part defined above does not exceed 50% of the left hand side of (2) to ensure sufficient sensitivity for the resonance part; (iii) the continuum threshold is determined by the requirement of maximum flatness of $m_\omega^2(n,\mathcal{M}^2,s_\omega)$ within the Borel window; and (iv) m_ω^2 follows as average with respect to $\mathcal{M}^2$.

Despite the linear density expansion of the condensates entering the coefficients c_j, the sum rule (4) is nonlinear in density. It is instructive to consider the linearized form. Using the notation $s_\omega = s_\omega^{(0)} + s_\omega^{(1)}n$ and $c_j = c_j^{(0)} + c_j^{(1)}n$ we arrive at

$$m_\omega^2(n,\mathcal{M},s_\omega^{(0)},s_\omega^{(1)}) = R + \Delta n \tag{5}$$

with

$$R = \frac{1}{N}\left\{c_0^{(0)}\mathcal{M}^2\left[1-\left(1+\frac{s_\omega^{(0)}}{\mathcal{M}^2}\right)E\right]-\frac{c_2^{(0)}}{\mathcal{M}^2}-\frac{c_3^{(0)}}{\mathcal{M}^4}-\frac{c_4^{(0)}}{2\mathcal{M}^6}\right\}, \tag{6}$$

$$N = c_0^{(0)}(1-E)+\frac{c_1^{(0)}}{\mathcal{M}^2}+\frac{c_2^{(0)}}{\mathcal{M}^4}+\frac{c_3^{(0)}}{2\mathcal{M}^6}+\frac{c_4^{(0)}}{6\mathcal{M}^8}, \tag{7}$$

$$\Delta = \frac{1}{N\mathcal{M}^2}\left\{\left[\frac{9R}{4M_N}+c_0^{(0)}Es_\omega^{(1)}(s_\omega^{(0)}-R)-c_2^{(1)}\left(1+\frac{R}{\mathcal{M}^2}\right)\right]\right.$$
$$\left.-\frac{c_3^{(1)}}{\mathcal{M}^2}\left(1+\frac{R}{2\mathcal{M}^2}\right)-\frac{c_4^{(1)}}{2\mathcal{M}^4}\left(1+\frac{R}{3\mathcal{M}^2}\right)\right\}, \tag{8}$$
$$E = e^{-s_\omega^{(0)}/\mathcal{M}^2},$$

which we use for illustrative purposes. The quantity R determines the vacuum properties of the ω meson; for the sake of estimates we can put it equal to $m_\omega^{(0)2}$ and use $\mathcal{M} \sim 1$ GeV. N contains only vacuum quantities, and $N > 0$ holds. Therefore, the sign of the in-medium shift of m_ω^2 is determined by Δ. For its estimate we note

$$c_2^{(1)} = \frac{1}{2}\left(1+\frac{1}{3}\frac{\alpha_s}{\pi}\right)\sigma_N-\frac{M_{N,0}}{27}+\left(\frac{1}{4}-\frac{5}{36}\frac{\alpha_s}{\pi}\right)A_2^{u+d}M_N$$
$$-\frac{9}{16}\frac{\alpha_s}{\pi}A_2^G M_N, \tag{9}$$

$$c_3^{(1)} = -\frac{112}{81}\pi\alpha_s\frac{\sigma_N\langle\bar{q}q\rangle_0}{m_q}\kappa_N-\left(\frac{5}{12}+\frac{67}{144}\frac{\alpha_s}{\pi}\right)A_4^{u+d}M_N^3$$
$$+\frac{615}{864}\frac{\alpha_s}{\pi}A_4^G M_N^3$$
$$+\frac{1}{4}M_N\left[\frac{3}{2}K_{u,1}+\frac{3}{8}K_{u,2}+\frac{15}{16}K_{u,g}\right]-\frac{7}{144}\sigma_N M_N^2, \tag{10}$$

where we include three active flavors on a 1 GeV scale; $M_{N,0} = 0.77$ GeV is the nucleon mass in the chiral limit. Some of the quantities in (6)–(10) are rather well-known (e.g., the twist-2 contributions), while others are individually less accurately fixed. We use for our evaluations $\langle\bar{q}q\rangle_0 = (-0.245\ \text{GeV})^3$, $\sigma_N = 0.045$ GeV, $\alpha_s = 0.38$, $Q_0 = 0.15$ GeV, $f_\pi = 0.093$ GeV, $A_2^{u+d} = 1.02$, $A_2^G = 0.83$, $A_4^{u+d} = 0.12$, $A_4^G = 0.04$, $K_{u,1} = -0.112\ \text{GeV}^2$, $K_{u,2} = 0.11\ \text{GeV}^2$, $K_{u,g} = -0.3\ \text{GeV}^2$ [11] to get $c_2^{(1)} \approx 0.17$ GeV, and $c_3^{(1)} \approx 0.2(\kappa_N - 0.7)\ \text{GeV}^3$ to obtain finally

$$\Delta \approx (4-\kappa_N)\frac{0.03}{n_0}\ \text{GeV}^2, \tag{11}$$

where we employ $s_\omega^{(0)} = 1.4\ \text{GeV}^2$, $s_\omega^{(1)} = -0.15 n_0^{-1}\ \text{GeV}^2$ (with $n_0 = 0.15\ \text{fm}^{-3}$ as nuclear saturation density) from an evaluation basing on (4) and neglect the c_4 term for the moment. To probe the uncertainty caused by less constrained quantities in (6)–(10) we assign N, R, $c_0^{(0)}$, $c_2^{(1)}$, and $c_3^{(1)}$ (the term in front of κ_N and the remainder separately) the large uncorrelated variations of $\pm 10\%$ and arrive at $\Delta = (2.8\cdots 5.3 - \kappa_N)\times(0.023\cdots 0.035)n_0^{-1}\ \text{GeV}^2$. In essence, to achieve a negative value of Δ and thus the experimentally observed [6] dropping of m_ω in medium, a sufficiently large value of κ_N is required, as evidenced by Eqs. (5) and (11). Thereby, the term $\propto c_3^{(1)}$ provides a counterbalance to the large Landau damping [17] [first term in (8)]. (For the ρ meson the Landau damping term is 9 times smaller [17], resulting in an always negative shift parameter conforming with the dropping ρ mass scenario in Ref. [18] and in qualitative agreement with the Brown-Rho scaling [19].)

Indeed, the evaluation of the complete sum rule (4) requires for the described parameter set $\kappa_N \geq 4$ to have $m_\omega^2(n > 0) < m_\omega^{(0)2}$; see Fig. 1. In other words, the above mentioned four-quark condensates must change drastically in the nuclear medium. With the above quoted parameters this translates into the huge amount of more than a 50% drop of the combined four-quark condensates at nuclear matter saturation density when relying on the linear density expansion up to such density. (The experiment [6] probes actually densities $\sim 0.6 n_0$.) Phrased differently, the density

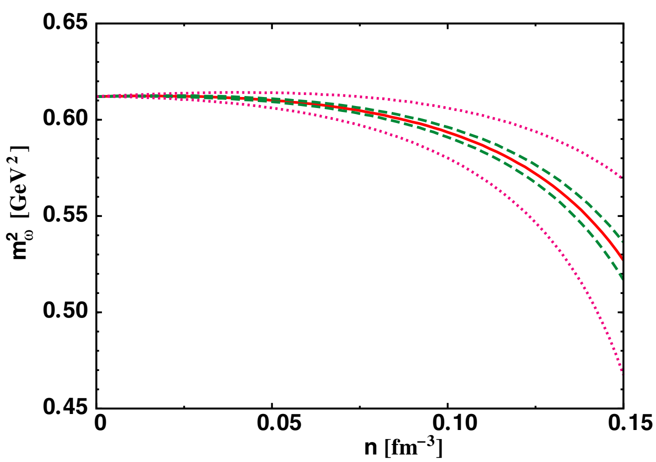

FIG. 1 (color online). The mass parameter m_ω^2 defined in Eq. (3) and averaged within the Borel window as a function of the baryon density for $\kappa_N = 4$ and $c_4 = 0$ (solid curve). Note that the parameter m_ω^2 coincides only in zero-width approximation with the ω pole mass squared; in general it is a normalized moment of $\text{Im}\Pi^\omega$ to be calculated from data or models. The sum rule Eq. (4) is evaluated as described in the text with appropriately adjusted κ_0. Inclusion of $c_4^{(0)} = \mathcal{O}(\pm 10^{-3})\ \text{GeV}^8$ requires a readjustment of κ_0 in the range $1\cdots 5$ to $m_\omega^{(0)2}$. A simultaneous change of κ_N in the order of 20% is needed to recover the same density dependence as given by the solid curve at small values of n. The effect of a $c_4^{(1)}$ term is exhibited, too ($c_4^{(1)} = \pm 10^{-5} n_0^{-1}\ \text{GeV}^8$, dashed curves; $c_4^{(1)} = \pm 5\times 10^{-5} n_0^{-1}\ \text{GeV}^8$, dotted curves; the upper (lower) curves are for negative (positive) signs.

dependence of the c_3 term must be stronger than the simple factorization allows.

To have some confidence in our estimate, the influence of c_4 must be evaluated. An order-of-magnitude estimate utilizing ground state saturation would yield $c_4 = (5.48 + 0.122n/n_0) \times 10^{-5}$ GeV^8 when considering only the first seven mass dimension-8 scalar condensates [20] and the two twist-2 condensates [21]. The corresponding value of $c_4^{(0)}$ is substantially smaller than the standard estimates related to the τ decay: [22] quotes $(-7 \cdots + 4) \times 10^{-3}$ GeV^8 pointing to some uncertainty also of $c_4^{(1)}$. We consider, therefore, c_4 as parameter and study its impact on the sum rule, as illustrated in Fig. 1.

Experimentally, a much stronger drop of the ω meson mass squared is found [6] than the in-medium change of m_ω^2 exhibited in Fig. 1. The conservative estimate of $\kappa_N > 4^{+0.7}_{-0.7} - 2.8^{+0.2}_{-0.4} \times 10^3$ GeV$^{-8} n_0 c_4^{(1)}$ obtained from an evaluation of the sum rule (4) expresses a condition for a decreasing value of m_ω^2 in the nuclear medium (the indicated variation of the numbers arise from an assumed uncertainty of $m_\omega^{(0)2}$ by $\pm 10\%$ which should reflect the approximate character of the vacuum sum rule). In other words, as long as $c_4^{(1)} < 1.7^{+0.3}_{-0.6} \times 10^{-3} n_0^{-1}$ GeV^8 a finite positive value of κ_N is required, i.e., a noticeable density dependence of the combined four-quark condensates. Note that this statement is independent of a model for $\mathrm{Im}\Pi^\omega$; it is based only on the observation that m_ω^2 must become smaller if additional strength of $\mathrm{Im}\Pi^\omega$ below $m_\omega^{(0)2}$ occurs in the medium, as observed in Ref. [6].

Finally, we mention that many more four-quark condensates enter other sum rules in different combinations. For instance, the investigation of the three coupled sum rule equations for the nucleon [23] points to some cancellations among the four-quark condensates when using the estimates from Ref. [24]. (The results of Ref. [24] cannot be employed directly for our ω sum rule since the flavor-mixing four-quark condensates are delivered in color combinations suitable only for the nucleon sum rule.) Furthermore, a crucial point is that the genuine chiral condensate is not suppressed in the nucleon sum rule, while in the ω sum rule it is.

In summary we argue that the recent CB-TAPS experiment [6] implies a noticeable drop (more than 50% when extrapolating to nuclear saturation density and truncating the sum rule beyond mass dimension 6) of a certain combination of four-quark condensates. Four-quark condensates are fundamental quantities, among others, characterizing the nonperturbative QCD vacuum. Specific four-quark condensates, changing under chiral transformation, represent further important order parameters for chiral symmetry restoration. Clearly, also other channels, besides the omega meson considered here, must be studied to gain more information on these particular four-quark condensates.

The authors thank R. Hofmann, S. Leupold, and W. Weise for stimulating discussions, and V. Metag and D. Trnka for informing us on the status of the CB-TAPS experiment. The work is supported by BMBF 06DR121 and GSI-FE.

[1] R. Rapp, J. Wambach, Adv. Nucl. Phys. **25**, 1 (2000).
[2] G. Colangelo, J. Gasser, and H. Leutwyler, Phys. Rev. Lett. **86**, 5008 (2001).
[3] S. Narison, *QCD as a Theory of Hadrons* (Cambridge University Press, Cambridge, England, 2004).
[4] M. Birse, Phys. Rev. C **53**, R2048 (1996).
[5] T. Doi, N. Ishii, M. Oka, and H. Suganuma, Phys. Rev. D **70**, 034510 (2004).
[6] D. Trnka *et al.* (CB-TAPS), Phys. Rev. Lett. **94**, 192303 (2005).
[7] M. Naruki *et al.*, nucl-ex/0504016.
[8] M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, Nucl. Phys. **B147**, 385 (1979).
[9] T. Schäfer and E. V. Shuryak, Phys. Rev. Lett. **86**, 3973 (2001).
[10] S. Leupold, W. Peters, and U. Mosel, Nucl. Phys. **A628**, 311 (1998).
[11] S. Zschocke and B. Kämpfer, Phys. Rev. C **70**, 035207 (2004).
[12] S. Narison and V. I. Zakharov, Phys. Lett. B **522**, 266 (2001).
[13] F. Klingl, N. Kaiser, and W. Weise, Nucl. Phys. **A624**, 527 (1997).
[14] S. Zschocke, B. Kämpfer, and O. P. Pavlenko, Eur. Phys. J. A **15**, 529 (2002); Phys. Lett. B **562**, 57 (2003); B. Kämpfer and S. Zschocke, Prog. Part. Nucl. Phys. **53**, 317 (2004).
[15] S. Leupold, Phys. Lett. B **616**, 203 (2005).
[16] P. Mühlich, T. Falter, and U. Mosel, Eur. Phys. J. A **20**, 499 (2004).
[17] A. K. Dutt-Mazumder, R. Hofmann, and M. Pospelov, Phys. Rev. C **63**, 015204 (2001).
[18] T. Hatsuda and S. H. Lee, Phys. Rev. C **46**, R34 (1992).
[19] G. E. Brown and M. Rho, Phys. Rep. **363**, 85 (2002) and further references therein.
[20] L. J. Reinders, H. Rubinstein, and S. Yazaki, Phys. Rep. **127**, 1 (1985).
[21] A. J. Buras, Rev. Mod. Phys. **52**, 199 (1980).
[22] S. Ciulli, C. Sebu, K. Schilcher, and H. Spiesberger, Phys. Lett. B **595**, 359 (2004).
[23] T. D. Cohen, R. J. Furnstahl, D. K. Kriegel, and X. M. Jin, Prog. Part. Nucl. Phys. **35**, 221 (1995).
[24] E. G. Drukarev, M. G. Ryskin, V. A. Sadovnikova, Th. Gutsche, and A. Faessler, Phys. Rev. C **69**, 065210 (2004).