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QCD sum rules for vector mesons in the nuclear medium

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A consistent treatment of quantum chromodynamical (QCD) sum rules in the nuclear medium is developed. Its close relation to the structure functions of the nucleon in the deep inelastic scattering is emphasized. The formalism is applied to the spectral changes of vector mesons (ρ, ω, and ϕ) in the nuclear matter. A linear decrease of the masses as a function of density is found. The four-quark condensate $\langle(\bar{q}q)^2\rangle$ and a twist-two condensate $\langle\bar{q}\gamma_\mu \bar{q}_\mu D_\nu q\rangle$ in medium play dominant roles for the mass-shift of light mesons. Physical implications of the result in finite nuclei and in heavy ion collisions are also discussed.

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The quantum chromodynamical (QCD) description of the elementary excitations in the nuclear medium is one of the current problems in hadron physics, since it has direct relevance to the phenomena in finite nuclei and in the relativistic heavy ion collisions [1,2]. In this Rapid Communication, we will report a first serious attempt to calculate the resonance parameters of the vector mesons (ρ, ω, and ϕ) in medium with nonzero baryon density (ρ) on the basis of the QCD sum rules (QSR). As we will show later, the system with $\rho\neq 0$ is the better place to study the partial restoration of chiral symmetry and associated phenomena than the finite temperature (T) system with $\rho=0$.

The purpose of this paper is twofold: The physical aim is to make a solid connection between QCD and the various suggestions related to the vector mesons in finite nuclei [2] and in the relativistic heavy ion collisions [3,4]. Another aim is to point out and clarify a vital role of the twist-two quark condensates in QSR at finite ρ which do not show up in the vacuum QSR [5] because of the Lorentz invariance. We will analyze such new terms by relating them to the structure functions of the deep-inelastic lepton-nucleon scattering (DIS).

In the following, the $N=Z$ cold nuclear matter with density ρ is taken as a ground state. We consider a relatively low density region, where the first-order expansion by ρ is valid and a model independent prediction is possible. (The validity range of this approximation will be discussed later.) Let us start with the longitudinal part of the current correlation function with zero three-momentum in the rest frame of matter:

$$\Pi_L(\omega^2)=\lim_{\mathbf{q}\to 0}[\Pi^\mu_\mu(\omega,\mathbf{q})/(-3\omega^2)]$$

with

$$\Pi_{\mu\nu}(\omega,\mathbf{q})=i\int d^4x\, e^{iqx}\langle TJ_\mu(x)J_\nu(0)\rangle_\rho\ , \tag{1}$$

where $J_\mu=(\bar{u}\gamma_\mu u\mp\bar{d}\gamma_\mu d)/2$ ($-$ for the ρ meson and $+$ for the ω meson) and $\bar{s}\gamma_\mu s$ (for the ϕ meson). $\langle\ \rangle_\rho$ denotes the ground-state expectation value, while we will use $\langle\ \rangle_0$ as a vacuum expectation value. The operator product expansion (OPE) of Eq. (1) gives

$$\mathrm{Re}\Pi_L(\omega^2)=\sum_n C_n(\omega^2,\mu^2)\langle\mathcal{O}_n(\mu^2)\rangle_\rho\ , \tag{2}$$

with μ being a renormalization point of the local operators $\mathcal{O}_n$. The Wilson coefficients are independent of density and all the medium effects are hidden in the expectation values $\langle\mathcal{O}_n\rangle_\rho$. Furthermore, the operators with uncontracted Lorentz indices can survive in the medium, since we have only O(3) invariance in the ground state. One such example is the vector operator $\bar{q}\gamma_\mu q$; its time component is nothing but the quark number density which is equal to 3ρ [6]. (This operator, however, does not show up in OPE for the vector mesons.)

By using the external field method with the Fock-Schwinger gauge [7], one can list all the possible local operators *with and without spin* up to six dimensions and calculate their Wilson coefficients. This procedure is similar to OPE in DIS [8] except that our expansion parameter is the dimension of $\mathcal{O}_n$ instead of twist. For example, the operators with nonzero Wilson coefficients in the chiral limit include [9] (i) four-quark operators with six dimensions such as $\bar{q}\Gamma_\mu\lambda^a q\cdot\bar{q}\Gamma_\nu\lambda^a q$, (ii) quark operators with four and six dimensions such as $\bar{q}\gamma_\mu D_\nu q$ and $\bar{q}\gamma_\mu D_\nu D_\lambda D_\rho q$ which are related to the twist-two operators in DIS, (iii) mixed operators like $\bar{q}D_\mu\tilde{G}_{\nu\alpha}\gamma^\alpha\gamma_5 q$, and (iv) pure gluon operators such as $G^2=-2(E^2-B^2)$, E^2+B^2, and $GDDG$. Note that (a) the tensor or scalar condensates of quark bilinear do not arise in the vector channel in the lowest order of α_s, and (b) only the operators with an odd number of covariant derivative (D_μ) appear among the twist-two operators. The detailed account will be given in Ref. [10].

Now let us examine the ρ dependence of the matrix elements of these operators. Up to linearity in ρ, one can always make a decomposition $\langle\mathcal{O}\rangle_\rho=\langle\mathcal{O}\rangle_0+\langle\mathcal{O}\rangle_N\rho$ where $\langle\ \rangle_N$ denotes the spin-averaged nucleon matrix element with the noncovariant normalization for the nucleon state $\langle p|p'\rangle=(2\pi)^3\delta^3(\mathbf{p}-\mathbf{p}')$. Since the Fermi motion of nucleons just increases the power of ρ, one can regard $\langle\ \rangle_N$ as a matrix element taken by the static nucleon ($\mathbf{p}=0$). The scalar condensates, which have both $\langle\mathscr{O}\rangle_0$ and $\langle\mathscr{O}\rangle_N$, can be written as (see, e.g., Ref. [11])

$$\langle\bar{u}u\rangle_\rho=\langle\bar{u}u\rangle_0+\frac{\Sigma_{\pi N}}{2\hat{m}}\rho\ ,$$
$$\langle\bar{s}s\rangle_\rho=\langle\bar{s}s\rangle_0+y\frac{\Sigma_{\pi N}}{2\hat{m}}\rho\ , \qquad (3)$$
$$\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle_\rho=\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle_0-\frac{8}{9}M_N^{(0)}\cdot\rho\ ,$$

where $\Sigma_{\pi N}\equiv\hat{m}\langle\bar{u}u+\bar{d}d\rangle_N$ is the π-$N\Sigma$ term with $\hat{m}$ being the averaged value of the current masses of u and d quarks, $y\equiv2\langle\bar{s}s\rangle_N/(\langle\bar{u}u\rangle_N+\langle\bar{d}d\rangle_N)$ is the strangeness content of the nucleon and $M_N^{(0)}$ is the nucleon mass in the chiral limit defined by the QCD trace anomaly. We have suppressed the obvious μ^2 dependence in the above relations. Empirical values of $\Sigma_{\pi N}$ and $\hat{m}$ (1 GeV) read (45±7) MeV [12] and(7±2) MeV [13], respectively. y and $M_N^{(0)}$ are not well determined yet and we quote two typical values here: $(y,M_N^{(0)})\sim(0.12,\ 830\ \text{MeV})$ (Nambu-Jona-Lasinio model [14]) and $(y,M_N^{(0)})\sim(0.22,\ 770\ \text{MeV})$ (improved chiral perturbation [12,15]). At normal nuclear matter density ($\rho_0=0.17\ \text{fm}^{-3}$), absolute values of the condensates decrease by 20–30 %, 4–8 % and ~8 % for $\langle\bar{u}u\rangle_\rho$, $\langle\bar{s}s\rangle_\rho$, and $\langle(\alpha_s/\pi)G^2\rangle_\rho$, respectively [16]. Such shifting of the condensates at finite ρ is much faster than those at finite T, since the density (temperature) correction to the condensates starts with linearity in ρ (quadratic in T).

For the four-quark condensates, by combining the linear ρ approximation and the ground-state saturation, i.e., the Hartree approximation (which is a similar assumption to the vacuum saturation in the vacuum QSR [5]), one can show that only the product of the scalar condensate is relevant, i.e., $\langle\bar{q}q\rangle_\rho^2$ expanded up to linear density $(=\langle\bar{q}q\rangle_0^2+2\langle\bar{q}q\rangle_0\langle\bar{q}q\rangle_N\rho)$.

The quark bilinear operators with four and six dimensions have nonvanishing matrix element only in medium. They can be related to the corresponding twist-two operators in DIS (and hence the parton distribution) aside from six-dimensional mixed condensates which we will discuss later:

$$\langle\mathscr{S}\bar{q}\gamma_{\mu_1}D_{\mu_2}\cdots D_{\mu_n}q(\mu^2)\rangle_N=(-i)^{n-1}A_n^q(\mu^2)T_{\mu_2\cdots\mu_n}\ , \qquad (4)$$
$$A_n^q(\mu^2)=2\int_0^1dx\,x^{n-1}[q(x,\mu^2)+(-)^n\bar{q}(x,\mu^2)]\ .$$

Here $\mathscr{S}$ makes the operator symmetric and traceless with respect to the Lorentz indices. $T_{\mu_1\cdots\mu_n}=[p_{\mu_1}\cdots p_{\mu_n}-$(trace terms)$]/2M_N$, with p_μ being the nucleon four-momentum ($p_\mu p^\mu=M_N^2$), and the factor $1/2M_N$ comes from our choice of the noncovariant normalization for the state vector. The "trace terms," which contain at least one $g_{\mu\nu}$ are as important as the leading term in contrast to the DIS case. $q(x,\mu^2)$ is the usual parton distri-

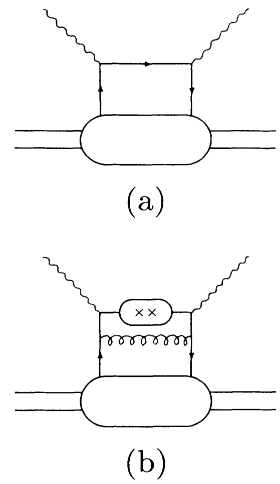

FIG. 1. (a) The contribution of $\langle\bar{q}\gamma_\mu D_\nu q\rangle_\rho$ to Π_L in Eq. (2). (b) One of the contributions of the four-quark operator to Π_L.

bution at scale μ^2. Since the relevant μ in QSR is the scale of the optimal Borel mass (~1 GeV), we evolve $q(x,\mu^2)$ from an initial distribution at $\mu^2\simeq10\ \text{GeV}^2$ to that at low μ^2. We adopt the LO scheme in Ref. [17] for this μ^2 evolution to calculate the nth moment $A_n^q(\mu^2)$. As μ^2 becomes small, the sea-quark and gluon components tend to be "absorbed" into the valence quark distribution as the perturbative QCD tells us.

As for the pure gluonic operators newly appearing in the medium (such as E^2+B^2 and $GDDG$), there is always a well-known quantum suppression in the Wilson coefficients compared to the quark operators of the same dimension [5]. Furthermore, the nucleon matrix elements of these operators, which is related to the moments of gluon distribution $g(x,\mu^2)$, tends to be suppressed at low μ^2 as we mentioned. Therefore we can safely neglect them compared to the quark contributions of the same dimension [18]. For example, the term proportional to E^2+B^2 in Eq. (2) turns out to be a few percent of that proportional to $\bar{q}\gamma_\mu D_\nu q$. As for the mixed condensate of six dimensions, some of them reduce to the four-quark condensate by the equation of motion which we have already considered. Others are supposed to be small, since the gluon component in the nucleon wave function is again small at low μ^2, which is explicitly discussed by Jaffe and Soldate [8] in the context of DIS. We will neglect them compared to the four-quark operators following their argument.

Summing up all these, one finds that the main ρ dependence in OPE originates from the four-quark condensate and $\langle\bar{u}\gamma_\mu D_n u\rangle_\rho$ for ρ and ω mesons and $m_s\langle\bar{s}s\rangle_\rho$ for the ϕ meson. In Fig. 1, the diagrammatic illustrations of the first two are shown. To get the resonance parameters, we use the standard dispersion relation in medium [19]

$$\Pi_L(\omega^2)=\frac{1}{\pi}\int_0^\infty du^2\frac{\text{Im}\Pi_L(u^2)}{(u-i\epsilon)^2-\omega^2}+\text{subtractions}\ . \qquad (5)$$

$\text{Im}\Pi_L(u^2)$ (hadronic part) is composed of three terms:

the scattering term, the resonance, and the continuum [20],

$$R\,\mathrm{Im}\Pi_L(u^2)=\delta(u^2)\rho_{sc}+F\delta(u^2-m^2)+\left[1+\frac{\alpha_s}{\pi}\right]\theta(u^2-S_0)\ , \quad (6)$$

with $R=4\pi$ (8π) for the ϕ meson (ρ and ω mesons) and ρ_{sc} denotes the contribution of the Landau damping on the hadronic side [4]. Matching the left- and right-hand sides of Eq. (5) in the asymptotic region $\omega^2\to-\infty$ by the method of Borel sum rule (BSR) [5] or the finite energy sum rule (FESR) [21], one can relate the resonance parameters to the ρ-dependent condensates. Here, for the qualitative argument, we will write down the FESR for the ρ (ω) meson in the chiral limit:

$$F-S_0\left[1+\frac{\alpha_s}{\pi}\right]=-2\pi^2 M_N^{-1}\cdot\rho\ ,$$

$$Fm^2-\frac{S_0^2}{2}\left[1+\frac{\alpha_s}{\pi}\right]=-Q_4-2\pi^2 A_2^{u+d}M_N\rho\ , \quad (7)$$

$$Fm^4-\frac{S_0^3}{3}\left[1+\frac{\alpha_s}{\pi}\right]=-Q_6-\frac{10}{3}\pi^2 A_4^{u+d}M_N^3\rho\ ,$$

where Q_4 and Q_6 are $\pi^2/3\langle(\alpha_s/\pi)G^2\rangle_\rho$ and $\frac{896}{81}\pi^3\langle\sqrt{\alpha_s}\bar{q}q\rangle_\rho^2$, respectively, taken up to linearity in ρ. The first sum rule (the local duality relation) essentially identifies the pole residue with the threshold parameter. The rhs of the first equation, which is a contribution from the hadronic scattering term ρ_{sc}, is negligible at low density. In the second sum rule, the twist-two quark condensate (the second term in the rhs with $A_2^{u+d}\simeq 0.90$ at 1 GeV) plays a significant role: its magnitude becomes comparable to the *vacuum* gluon condensate with the same sign, which makes the vector meson mass smaller in the medium. (Remember that the gluon condensate contributes to the mass negatively.) The four-quark condensate in the third sum rule decreases by the medium effect according to Eq. (3), which also makes the vector meson mass smaller. (Remember that the four-quark condensate contributes to the mass positively.) Contribution of the twist-two condensate in the third sum rule is small because it corresponds to the higher moment of the structure function ($A_4^{u+d}\simeq 0.12$ at 1 GeV).

One can also write down similar sum rules for the ϕ meson. In this case, because of the OZI suppression, the contributions of twist-two operators are rather small ($A_2^s\simeq 0.05$ and $A_4^s\simeq 0.002$ at 1 GeV). However, still a small mass shift occurs mainly due to the decrease $m_s\langle\bar{s}s\rangle_\rho$ which is the main term in OPE for the ϕ meson.

To find the validity range of the linear ρ approximation for $\langle\mathcal{O}_n\rangle_\rho$, we have estimated some of the $\rho^{n>1}$ terms in the quark condensates. They include the $\rho^{5/3}$ term and the ρ^2 term in $\langle\bar{q}q\rangle_\rho$ due to the Fermi motion of nucleons and the two-body nucleon-nucleon interactions respectively, and also the ρ^2 terms that appear in the four-quark condensates $\langle(\bar{q}q)^2\rangle_\rho$ and $\langle(q^\dagger q)^2\rangle_\rho$. They affect at most 15% (30%) of the linear terms in ρ at $\rho=\rho_0$ ($2\rho_0$). Therefore, we take $0<\rho<2\rho_0$ as an optimistic validity range of our analysis. More systematic analysis of the higher order terms will be given in Ref. [10].

The numerical calculations for the density dependent masses shown in Fig. 2 have been done by using the BSR. We took one of the canonical procedures to determine the resonance parameters: First S_0 is determined to make the meson mass m least sensitive to the Borel mass M in the region $M_{\min}<M<M_{\max}$ determined by the 30% criterion. Then the minimum of $m(M)$ is adopted as the physical mass. We have also checked that the FESR has the same qualitative result with BSR. Since we are only interested in the medium modifications of the resonances, the same *vacuum* parameters with those in Ref. [22] are adopted without any fine tuning [23]; for instance, $Q_4=0.05$ GeV^4, $Q_6=0.06$ GeV^6, $m_s\langle\bar{s}s\rangle_0=-(210$ MeV$)^4$, etc. As for the parameters in Eq. (3), the central values are taken to calculate the curves in Fig. 2. The masses of ρ, ω, and ϕ mesons decrease almost linearly in density. (The deviation from the linearity above ρ_0 in the ρ-ω channel should be regarded as an uncertainty of our procedure of the Borel analysis [24].)

From Fig. 2, one can make a linear fit of $m(\rho)$ using

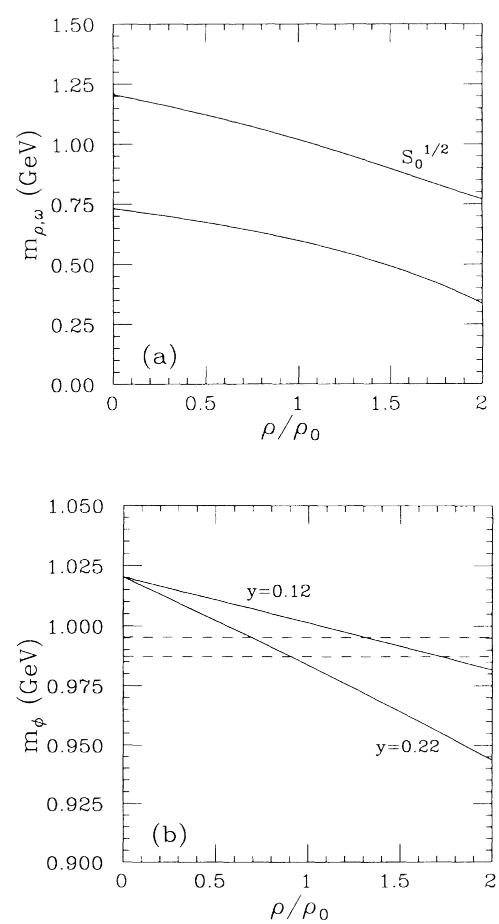

FIG. 2. (a) The ρ-ω meson mass m and the continuum threshold $S_0^{1/2}$ as a function of ρ/ρ_0. (b) The same figure for the ϕ-meson mass with two typical values of y (the strangeness content in the nucleon). Dashed lines indicate the K^0K^0 and K^+K^- threshold at $\rho=0$ which are the main decay modes of ϕ.

the values at $\rho=0$ and ρ_0, which results in

$$m(\rho)/m(0) \simeq 1 - C(\rho/\rho_0) , \quad (8)$$

with $C \simeq 0.18$ $(0.15y)$ for ρ-ω (ϕ). Error for the numbers 0.18 and 0.15 is $\sim \pm 30\%$ which comes from the uncertainty of $\Sigma_{\pi N}/2\hat{m}(1 \text{ GeV}) = 3.21 \pm 1.05$. In the ϕ channel, there is an extra uncertainty due to y which is shown in Fig. 2. The threshold parameter $\sqrt{S_0}$ has also a linear decrease. The above result tells us that vector meson masses almost scale like the quark condensate $\langle \bar{q}q \rangle_\rho$.

Much implicit evidence of the dropping ρ-ω mass in finite nuclei is extensively discussed by Brown and Rho in Ref. [2], which includes the K^+ scattering off medium-heavy nuclei, polarization observables in nucleon-nucleus scattering, the enhancement of the spin-orbit part of the G-matrix element, and the enhancement of tensor-force due to the ρ meson in nuclei. These phenomenological studies require $C \sim 0.15$ which is comparable to our value in Eq. (8). In this sense, our result seems to support the above phenomenological idea. However, one should also notice that $m(\rho)$ in our calculation is the pole position of the propagator in the medium, while the relevant quantity in most of the nuclear physics application is the "mass" defined in the space-like momentum (i.e., the screening mass). These two are different in general and one needs to study QSR with $\mathbf{q} \neq 0$ in Eq. (1) to check whether they are close or not. On the other hand, the direct measurement of the pole position in the medium is possible through the leptonic decay of vector mesons created in the hot/dense hadronic system by the relativistic heavy ion collisions [3,4] or created in finite nuclei by the electron machines. In particular, only a few percent change of the ϕ-meson mass can affect its decay property drastically, since m_ϕ is very close to the threshold of its main decay mode ($K^0\bar{K}^0$ and K^+K^-) [3]. Figure 2 suggests that the ϕ-meson is going to be more stable in medium, *if* the kaon mass (m_K) does not change in medium [25].

Now let us briefly compare our result at $(T=0, \rho \neq 0)$ with that at $T \neq 0, \rho=0$. In the latter case, the FESR at finite T tells $m(T)/m(0) \sim \langle \bar{q}q \rangle_T^{1/3} \sim (1-T^2/T_c^2)^{1/6}$ with $T_c \simeq 2f_\pi$ [1]. (Here we have assumed a mean field behavior for the condensate $\langle \bar{q}q \rangle_T \sim \sqrt{1-T^2/T_c^2}$.) This shows one can see the considerable change of $m(T)$ only near the critical point [26]. Therefore, the finite ρ system (such as the finite nuclei or the high density region of the heavy ion collisions) will be a better place to see the effect of the partial restoration of chiral symmetry and its associted changes of hadron properties in experiment.

To go beyond the linear density approximation, we have to introduce models for the behavior of the condensates at high ρ. This point together with the full details of the present paper and the discussions on the other vector mesons (A_1, K^*, and J/Ψ) will be reported elsewhere [10].

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