Fate of QCD sum rules or fate of vector meson dominance in a nuclear medium

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A current-current correlator with the quantum numbers of the omega meson is studied in a nuclear medium. Using weighted finite energy sum rules and dispersion relations for the current-nucleon forward scattering amplitude it is shown that strict vector meson dominance and QCD sum rules are incompatible with each other. This implies that at least one of these concepts — which are both very powerful in the vacuum — has to fade in the nuclear environment.

Keywords: QCD sum rules, vector meson dominance, meson properties, nuclear matter

I. INTRODUCTION

The present work addresses the interrelation of two important concepts of hadron physics, namely QCD sum rules [1] and vector meson dominance (VMD) [2]. Both concepts are extremely successful in the description of hadrons and their interactions. However, if we turn from the elementary (vacuum) properties and cross sections of hadrons to a many-body system (here: nuclear matter), we will demonstrate that the two concepts are not compatible any more. We will restrict ourselves here to the omega meson and comment briefly on possible extensions in the end. It is important to note that both methods are also frequently used for in-medium calculations — concerning QCD sum rules cf. e.g. [3, 4, 5, 6] and references therein, concerning VMD cf. e.g. [3, 7] and references therein. Our work shows that at least one of the two methods must be modified (at least for nuclear matter calculations and at least for the omega meson). In fact, recently the incompatibility of the two concepts has been observed numerically in [8]. Here we study that further by presenting an analytical derivation.

In the next section we will discuss the QCD sum rule method for our case of interest. In section III we introduce the concept of VMD and combine it with the previously derived sum rule formula. We will end up with an inconsistent relation. This proves that (at least) one of the input assumptions — sum rules or VMD — must be wrong for the studied case, omega mesons in nuclear matter. Finally we discuss this issue further in section IV.

II. QCD SUM RULES

In this work we study the properties of a vector-isoscalar current

\[ j_\mu := \frac{1}{2} (\bar{u} \gamma_\mu u + \bar{d} \gamma_\mu d) \]  

which is at rest with respect to the nuclear medium. From the current-current correlator

\[ \Pi_{\mu\nu}(q) = i \int d^4x e^{iqx} \langle T j_\mu(x) j_\nu(0) \rangle_{\text{med}} \]  

we construct the dimensionless quantity

\[ R(q^2) := \frac{\Pi^\rho_{\mu\nu}(q^2, q^2 = 0)}{-3q^2}. \]  

\( R \) has a direct physical meaning in the time-like region \( s = q^2 > 0 \). It is related to the cross section \( e^+e^- \rightarrow \text{hadrons} \) with isospin 0 via \( \bar{2} \)

\[ \frac{\sigma_I=0(e^+e^- \rightarrow \text{hadrons})}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} = \frac{4}{3} \pi \text{Im}R. \]  

At least for low energies this cross section is determined by hadronic degrees of freedom. For high energies the cross section can be calculated by perturbative QCD. One supposes that perturbative QCD yields a reliable result for energies above the so-called continuum threshold \( \bar{1} \). One gets

\[ \text{Im}R(s) = \frac{1}{8\pi} \left( 1 + \frac{\alpha_s}{\pi} \right) \quad \text{for} \quad s > s_0(\rho_N). \]  

\( \bar{3} \)

\[ \bar{4} \]
The continuum threshold $s_0$ might depend on the nuclear density $\rho_N$, but otherwise the high-energy form is assumed to be independent of the medium. This is in agreement with the general picture that in-medium changes are a (collective) long-distance effect.

As outlined e.g. in \[9\] in-medium QCD sum rules can be obtained from an off-shell dispersion relation which integrates over the energy at fixed (here vanishing) three-momentum of the current. We also restrict ourselves to small densities $\rho_N$ by using the linear-density approximation. Effectively this means that the current is at rest with respect to the nucleon on which it scatters. The lowest two finite energy sum rules (FESRs) are given by \[10, 11, 12\]

\[
\frac{1}{\pi} \int_0^{s_0(\rho_N)} ds \text{ Im} R(s, \rho_N) = c_0 s_0(\rho_N) - \frac{9\rho_N}{4m_N}, \quad (6a)
\]

\[
\frac{1}{\pi} \int_0^{s_0(\rho_N)} ds s \text{ Im} R(s, \rho_N) = \frac{1}{2} c_0 s_0^2(\rho_N) - c_2. \quad (6b)
\]

Here $\rho_N$ denotes the nuclear density and the $c_i$ encode the condensates:

\[
c_0 = \frac{1}{8\pi^2} \left( 1 + \frac{\alpha_s}{\pi} \right), \quad (7a)
\]

\[
c_2 = m_q \langle \bar{q}q \rangle_{\text{med}} + \frac{1}{24} \left( \frac{\alpha_s}{\pi} G^2 \right)_{\text{med}} + \frac{1}{4} m_N a_2 \rho_N. \quad (7b)
\]

Note that we have followed the common practice to neglect contributions from $\alpha_s$ suppressed twist-two operators (cf. e.g. \[13\] for details). Their inclusion would not change our lines of reasoning. The quantity $a_2$ is a moment of parton distributions and therefore well constrained by deep inelastic scattering \[4\]. Note that we have extracted from $R$ the Landau damping contribution — last term on the right hand side of \[6a\].

Using the linear-density approximation we get for quark and gluon condensate \[4\]:

\[
m_q \langle \bar{q}q \rangle_{\text{med}} = m_q \langle \bar{q}q \rangle_{\text{vac}} + m_q \langle N|\bar{q}q|N\rangle \rho_N = -\frac{1}{2} F_p^2 M^2 + \frac{1}{2} \sigma_N \rho_N, \quad (8)
\]

and

\[
\left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle_{\text{med}} = \left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle_{\text{vac}} - \frac{8}{11 - \frac{4}{3}N_f} m_N^{(0)} \rho_N. \quad (9)
\]

Let us briefly discuss the essence of the sum rules \[8\]: The basic idea \[1\] is that the hadronic low-energy spectral information (which enters the left hand side) is connected to QCD information about the non-perturbative vacuum, or in other words: to perturbative QCD improved by condensates via an operator product expansion. In short: High-energy information (about quarks and gluons) is connected to low-energy information (about hadrons). As already discussed in \[1\], it is not always possible to connect these informations. At least, one has to make sure that the hadronic side of a sum rule is indeed given/dominated by the low-energy information and that the condensate side is dominated by the high-energy information encoded in the known condensates of low dimensionality. We will come back to that point below.

Besides the condensates which are (more or less) well known and the quantity $\text{Im} R$ we are interested in, we also have the continuum threshold $s_0$ and its density dependence. It introduces an additional parameter which we are not primarily interested in. In addition, the sum rules \[8\] mix vacuum and in-medium information while we are only interested in the latter. At least, we want to make sure that uncertainties in the vacuum description do not influence our conclusions for the in-medium changes. In the following, we will show that one can get better sum rules which solve some of the mentioned problems.

To be more sensitive to the in-medium modifications we differentiate the FESRs \[8\] with respect to the density.

\[\text{Sign and size of the Landau damping contribution are discussed in some detail in } \[8\].\]
Since we work in the linear-density approximation we should set \( \rho_N \) to zero after the differentiation. We get

\[
\frac{1}{\pi} \int_0^{s_0} ds \frac{\partial}{\partial \rho_N} \text{Im} R(s, \rho_N) \bigg|_{\rho_N=0} + \hat{c} = c_0 \frac{ds_0}{d\rho_N} \bigg|_{\rho_N=0} - \frac{9}{4m_N}, \tag{10a}
\]

\[
\frac{1}{\pi} \int_0^{s_0} ds \frac{\partial}{\partial \rho_N} \text{Im} R(s, \rho_N) \bigg|_{\rho_N=0} + s_0 \hat{c} = c_0 s_0 \frac{ds_0}{d\rho_N} \bigg|_{\rho_N=0} - \frac{dc_2}{d\rho_N}, \tag{10b}
\]

with \( s_0 = s_0(\rho_N = 0) \) and

\[
\hat{c} = \frac{1}{\pi} \text{Im} R(s_0, \rho_N) \frac{ds_0}{d\rho_N} \bigg|_{\rho_N=0} . \tag{11}
\]

The FESRs \( 6 \) — and also \( 10 \) — show an unpleasant feature: The region around the continuum threshold can sizably contribute to the respective integrals. On the other hand, we recall that the sum rules are derived under the assumption that at the continuum threshold perturbative QCD immediately sets in and gives a reliable description. This is expressed in \( 8 \). Clearly this is an over-simplifying assumption. Therefore, one would prefer sum rules which are only sensitive to the low-energy region one is interested in and insensitive to the region around the threshold. This criterion is not met by the FESRs, but e.g. by the Borel sum rules already used in the seminal QCD sum rule papers \( 1 \). However, also Borel sum rules have their shortcomings as compared to FESRs: In principle, condensates of arbitrary high dimensions enter the Borel sum rules while only condensates of a specific dimension enter a particular FESR (e.g. only dimension-four condensates collected in \( c_2 \) enter the sum rule \( 9 \)). In addition, Borel sum rules introduce additional parameters (the Borel window) in which one is primarily not interested in. A method which combines the advantages of FESR and Borel sum rules are weighted FESRs. For our case of interest we obtain the following weighted finite energy sum rule (WFESR) by a linear combination of the equations \( 10 \):

\[
\frac{1}{\pi} \int_0^{s_0} ds (s_0 - s) \frac{\partial}{\partial \rho_N} \text{Im} R(s, \rho_N) \bigg|_{\rho_N=0} = -\frac{9s_0}{4m_N} + \frac{dc_2}{d\rho_N}. \tag{12}
\]

This is the sum rule which we will use from now on.

As already noted, the standard FESRs are rather sensitive to the modeling of the transition region from the hadronic part to the continuum (see also e.g. \( 14 \) and references therein). This is different for \( 12 \), since the transition region is suppressed by \( (s_0 - s) \). Therefore \( 12 \) is more reliable as it is insensitive to details of the threshold modeling at \( s_0 \). This is a frequently discussed line of reasoning which is not special for in-medium sum rules (cf. e.g. \( 12 \) and references therein).\(^2\) The in-medium WFESR \( 12 \), however, has an additional feature: It is independent of the in-medium change of the threshold parameter, i.e. independent of \( \frac{ds_0}{d\rho_N} \). Note that the vacuum threshold \( s_0 \) appears in \( 12 \). This, however, can be fixed by an independent vacuum sum rule analysis which is free of all in-medium uncertainties. If one takes the arithmetic average of the squared masses of the omega and of its first excitation \( 17 \) one gets \( s_0 \approx 1.3 \text{ GeV}^2 \). We note already here that for our qualitative arguments we do not need the numerical value.

We would like to stress again that the sum rule \( 12 \) constitutes a big step forward in the sum rule analysis of in-medium properties: First, we are directly sensitive to in-medium changes in contrast to traditional analyses \( 3 \) \( 4 \) \( 5 \) \( 18 \) which study vacuum plus medium contributions. Second, we have got rid of all additional parameters like the in-medium continuum threshold and the Borel masses (cf. also the discussion in \( 11 \) \( 19 \)). Third, we still share with Borel type sum rules the feature that we are less sensitive to the modeling of the continuum threshold.

### III. VECTOR MESON DOMINANCE

Strict VMD implies the identification \( 2 \) \( 7 \)

\[
\hat{j}_\mu = -\frac{M^2}{g} \omega_\mu . \tag{13}
\]

\(^2\) Indeed, a very interesting and successful application of WFESRs and a failure of standard FESRs are documented in \( 16 \) for Weinberg type sum rules in the vacuum.
with the mass of the omega $M_\omega \approx 782\text{ MeV}$ \cite{17}, the universal vector meson coupling $g \approx 6$ \cite{18} and the omega meson field $\omega_\mu$. In other words, all hadronic interaction of the current \cite{11} with hadrons is mediated by the omega meson. One gets

$$\text{Im} R(s) = -\frac{M_\omega^4}{g^2 s} \frac{1}{s - M_\omega^2 - \Pi_\omega(s)}$$ \hspace{1cm} (14)

with the self energy $\Pi_\omega$, which in general splits into a vacuum and an in-medium part. Let us stress again, that according to strict VMD all in-medium modifications influence the propagator of the omega meson (via the self energy), but not the vertex between the current \cite{11} and the omega. In linear-density approximation we get

$$\text{Im} R(s, \rho_N) = -\frac{M_\omega^4}{g^2 s} \frac{1}{s - M_\omega^2 - \rho_N T(s)}$$ \hspace{1cm} (15)

and therefore

$$\frac{\partial}{\partial \rho_N} \text{Im} R(s, \rho_N) \bigg|_{\rho_N=0} = \frac{M_\omega^4}{g^2 s} \text{Im} \left[ \left( \frac{d}{ds} \frac{1}{s - M_\omega^2 + i\epsilon} \right) T(s) \right]$$ \hspace{1cm} (16)

with the omega-nucleon forward scattering amplitude $T(s)$ for an in general off-shell omega with invariant mass squared $s$ which is at rest with respect to the nucleon. Note that we have neglected the small vacuum self energy of the omega meson. We will come back to this approximation below.

Using \cite{10} we obtain for the left hand side of the sum rule \cite{12} after some algebra:

$$\frac{1}{\pi} \int_{0}^{s_0} ds \ (s_0 - s) \frac{\partial}{\partial \rho_N} \text{Im} R(s, \rho_N) \bigg|_{\rho_N=0} = \frac{M_\omega^4}{g^2 \pi} \left\{ \int_{0}^{s_0} ds \left[ s_0 \frac{\text{Re} T(s) + 1}{s - M_\omega^2 + i\epsilon} \text{Im} T(s) - \left( \frac{s_0}{s} - 1 \right) \frac{\text{Re} T(s)}{s - M_\omega^2 + i\epsilon} \text{Im} T'(s) \right] \right. \right.$$

$$\left. - \frac{s_0}{M_\omega^2} \pi \text{Re} T(M_\omega^2) + \left( \frac{s_0}{M_\omega^2} - 1 \right) \pi \text{Re} T'(M_\omega^2) \right\} .$$ \hspace{1cm} (17)

It is important to note that the integrals which appear in the last expression look like dispersive integrals. Indeed, we will use in the following dispersion relations to calculate $\text{Re} T$ and bring the non-integral terms on the right hand side of \cite{17} in a form similar to the integral terms. There is, however, one notable difference: Dispersive integrals cover the whole energy range, i.e. they are not restricted by $s_0$.

Next we relate the real part of the scattering amplitude to its imaginary part using a one time subtracted dispersion relation, i.e. \cite{212}

$$\text{Re} T'(s) = -\frac{1}{\pi} \text{P} \int_{0}^{\infty} ds' \frac{\text{Im} T'(s')}{s - s'} = -\frac{1}{\pi} \int_{0}^{\infty} ds' \frac{\text{Im} T'(s') \text{Re} \frac{1}{s - s' - i\epsilon}}{s}$$ \hspace{1cm} (18)

and therefore

$$\text{Re} T(s) = \text{Re} T(0) - \frac{s}{\pi} \int_{0}^{\infty} ds' \frac{\text{Im} T(s')}{s'} \text{Re} \frac{1}{s - s' - i\epsilon} .$$ \hspace{1cm} (19)

The subtraction constant appearing in \cite{19} is the omega-nucleon forward scattering amplitude at the photon point. VMD relates this quantity to the isospin-0 part of the photon-nucleon scattering amplitude for vanishing photon momenta. The latter is just Thomson scattering and therefore we get

$$\text{Re} T(0) = \frac{9g^2}{4m_N} .$$ \hspace{1cm} (20)

Now we use \cite{17}, \cite{18} and \cite{19} to evaluate the left hand side of the WFESR \cite{12}. We get after some lengthy but straightforward algebra

$$-\frac{s_0}{g^2} \text{Re} T(0) + \frac{1}{\pi} \frac{M_\omega^4}{g^2} X = -\frac{9s_0}{4m_N} + \frac{dc_2}{d\rho_N} .$$ \hspace{1cm} (21)
with

$$X = \frac{1}{M_\omega^2} \text{Im} T(s_0) - \frac{s_0}{M_\omega^2} \int_{s_0}^\infty ds \frac{\text{Im} T(s)}{s} \text{Re} \left( \frac{1}{s - M_\omega^2 + i\epsilon} \right) + \left( \frac{s_0}{M_\omega^2} - 1 \right) \int_{s_0}^\infty ds \text{Im} T'(s) \text{Re} \left( \frac{1}{s - M_\omega^2 + i\epsilon} \right).$$

(22)

Using finally \cite{24} we get

$$\frac{1}{\pi} \frac{M_\omega^4}{g^2} X = \frac{dc_2}{d\rho_N}. \quad (23)$$

It is important to note that $X$ contains purely high-energy information, i.e. the imaginary part of the scattering amplitude above the continuum threshold. Concerning the sum rule philosophy this is a disastrous result: The condensates on the right hand side — to be more precise: their in-medium changes — do not constrain the low-energy information as they should, but rather the high-energy information. In contrast, the basic idea of the sum rule method is to make contact between the condensates and the hadronic low-energy information. As already discussed above, the high-energy part (and especially the part around the continuum threshold) is much less accurately treated. In turn this means that a sum rule becomes unreliable, if it is not the low-energy information which is connected to the condensates, but part of the high-energy information. The sum rule method breaks down.

In the spirit of the sum rule method one might even take a somewhat more simplified point of view to demonstrate what is wrong with \cite{24}: We recall the proposition that the high-energy part of $\text{Im} R$ is unchanged by the medium, as expressed in \cite{19}. Consequently, it should not matter much whether one stops the dispersion integral \cite{19} at $s_0$ or at infinity. Stopping the dispersion at $s_0$, however, is equivalent to completely neglecting $X$. Hence, we would deduce from \cite{23} that the in-medium changes of the condensates vanish. On the other hand, plugging in reasonable numbers \cite{4, 8}, one finds that $dc_2/d\rho_N$ does not vanish. This demonstrates even more clearly the breakdown of the sum rule method.

\section*{IV. FURTHER DISCUSSION}

We have found that the sum rule method becomes unreliable for the in-medium part of the correlator. Note, however, that this result is caused by VMD. Without VMD there would be terms remaining at the left hand side of the sum rules which contain low-energy information. This could easily be checked e.g. with the extension used in \cite{21}. We will not go through this exercise here. Hence, the main result of our purely analytical calculations is the following: If strict VMD still works well in a nuclear medium, then the sum rule method does not work any more. In turn, this means, that if the sum rule method works in a nuclear medium, then strict VMD has to fade. We would like to stress that neither the sum rule method nor VMD is strictly derivable from QCD. Therefore, we will not speculate in the present paper which of the concepts should be modified. We only report our finding that they do not fit together for an in-medium situation. We note in passing that an in-medium fate of VMD has been found in the context of the hidden local symmetry approach \cite{21}.

We should discuss the methods and approximations used to obtain our result: By replacing \cite{18} by \cite{15} we have neglected the vacuum self energy of the omega meson, i.e. its width and the corresponding real part. Indeed, without that approximation the integrals in \cite{17} would not resemble dispersions. This would induce changes on the left hand side of \cite{23} which scale with the (vacuum) width of the omega meson. Do they contain the proper low-energy information which is constrained by condensates on the right hand side of \cite{23}? In principle, this could be a way to reconcile sum rules and VMD. We give two reasons why this cannot be the case: From a phenomenological point of view, we observe that the width of the omega is very small — about 1% of its mass \cite{17}. Such small modifications cannot account for the condensates. Also from a more formal point of view this would be unsatisfying: One can study the sum rules as a function of the number of colors $N_c$ and explore the large-$N_c$ behavior \cite{22, 23}. In general, the condensates \cite{7} are $O(N_c)$. This is also true for their in-medium changes.\footnote{The change of the gluon condensate is an exception, see e.g. \cite{24} for details.} Also the hadronic left hand side of the sum rules is $O(N_c)$ which can be seen e.g. from \cite{15} by noting that the omega mass and the scattering amplitude $T$ are $O(N_c^0)$ and the coupling $g^2 = \alpha(1/N_c)$. The vacuum width of the omega, however, is $O(1/N_c)$. Therefore, the inclusion of the vacuum width of the omega on the hadronic side of the sum rule would induce terms which are $O(N_c^0)$ or lower. This does not match with the condensate side. We conclude that the inclusion of a vacuum self energy for the omega does not change our lines of reasoning. We note in passing, that the large-$N_c$ arguments could be
immediately taken over for the rho meson. Neglecting the vacuum width of the rho meson one would also conclude that for a rho meson in a nuclear medium strict VMD does not fit together with the sum rules. On the other hand, from a phenomenological point of view it is less satisfying to neglect the rather large width \( (\approx 150 \text{ MeV}) \) of the rho meson. Therefore we have decided to concentrate on the omega meson in the present work.

One might ask how sensitive our arguments are concerning the choice of the used type of sum rules (Borel vs. FESRs etc.). We have already motivated why we prefer the use of WFESRs. On top of the arguments already given, we would like to stress that using a WFESR we were able to show the incompatibility of sum rules and VMD in a purely analytical way. On the other hand, the same incompatibility can also be found numerically within analyses using other types of sum rules — and other types of vector mesons. The sum rule practitioners have just not payed much attention to it. As an example we discuss the Borel sum rule analysis of [5] for the rho meson: The parameter \( F \) introduced in equation 12 in [5] should not change if strict VMD holds. From table 2 of [5] we find that \( F \) drops sizably in a medium for all parameter choices. We conclude that our findings are not special to the chosen type of sum rules. But obviously the proof is most elegant using a WFESR.

Finally, let us discuss in more detail which kind of VMD is actually incompatible with the in-medium sum rules. The relation (13) which holds on the level of currents and fields can be obtained from a Lagrangian

\[
\mathcal{L}_{\text{int}} = -\frac{e M_a^2}{3g} \omega^\mu A_\mu + \omega^\mu j_{\text{had}}^\mu
\]  

with the photon field \( A_\mu \), the electromagnetic coupling \( e \) and the hadronic current \( j_{\text{had}}^\mu \). Such a Lagrangian — albeit frequently used — is somewhat unsatisfying as it mixes photon and vector meson fields. In that way, the field \( A_\mu \) gets a mass and the massless “real” photon emerges from a linear combination of \( A_\mu \) and the vector meson field. Such complications are avoided by a Lagrangian

\[
\mathcal{L}_{\text{int}} = -\frac{e}{6g} \omega^{\mu \nu} F_{\mu \nu} + \omega^\mu j_{\text{had}1}^\mu + e A_\mu j_{\text{had}2}^\mu
\]  

with the field strengths \( \omega^{\mu \nu} \) and \( F_{\mu \nu} \) for vector meson and photon, respectively. Now the \( A_\mu \) field remains massless, but real photons decouple from the vector mesons. Therefore, a direct coupling of photons to hadrons is needed. In general, the Lagrangian (24) leads to a form for the current-current correlator which is more complicated than the one given in (13). In such a case, one can find in-medium interactions which do not contradict the sum rule setting. However, if there is a relation between the two currents which appear in (24), namely

\[
9 j_{\text{had}1}^\mu = j_{\text{had}2}^\mu
\]  

then relation (13) still holds. Since our proof of incompatibility relies on (13) and not on the strict VMD relation (13) we can conclude that the sum rules are also incompatible with an in-medium interaction generated from an extended VMD Lagrangian (25), if condition (26) holds.

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