Calculation of tensor susceptibility beyond rainbow-ladder approximation of Dyson-Schwinger equations approach

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Abstract

In this paper, we extend the calculation of tensor vacuum susceptibility in the rainbow-ladder approximation of the Dyson-Schwinger (DS) approach in [Y.M. Shi, K.P. Wu, W.M. Sun, H.S. Zong, J.L. Ping, Phys. Lett. B 639, 248 (2006)] to that of employing the Ball-Chiu (BC) vertex. The dressing effect of the quark-gluon vertex on the tensor vacuum susceptibility is investigated. Our results show that compared with its rainbow-ladder approximation value, the tensor vacuum susceptibility obtained in the BC vertex approximation is reduced by about 10%. This shows that the dressing effect of the quark-gluon vertex is not large in the calculation of the tensor vacuum susceptibility in the DS approach.

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The QCD vacuum susceptibilities play an important role in characterizing the non-perturbative aspects of QCD and in the determination of hadron properties \[1–3\]. In particular, tensor vacuum susceptibility is relevant for the determination of the tensor charge of the nucleon in the QCD sum rule approach \[4\]. The previous calculations of the tensor vacuum susceptibility have shown that the theoretical treatment of this quantity is subtle and different treatments can lead to different results \[5–9\]. In order to get a reliable theoretical prediction of the tensor charge, one needs to determine the tensor vacuum susceptibility as precisely as possible. Recently, a particular implementation of the vacuum polarization definition of the vector vacuum susceptibility has been proposed in Ref. \[10\], in which the vector vacuum susceptibility is totally determined by the dressed quark propagator and the dressed vector vertex. Soon this definition of vector vacuum susceptibility has been generalized to calculate the tensor vacuum susceptibility by some of the same authors in Ref. \[11\]. Just as was shown in Ref. \[11\], in order to calculate the tensor vacuum susceptibility, one needs to know the dressed quark propagator and the tensor vertex in advance. At present it is impossible to solve for the dressed quark propagator and the tensor vertex from first principles of QCD. So one has to resort to various nonperturbative QCD models. In the past few years, considerable progress has been made in the framework of the rainbow-ladder approximation of the Dyson-Schwinger (DS) approach \[12–17\]. Due to the great success of the rainbow-ladder approximation of the DS approach in hadron physics, the authors in Ref. \[11\] adopt this approximation to solve for the dressed quark propagator and the tensor vertex and from these obtain the numerical value of the tensor vacuum susceptibility. However, it is well known that the rainbow-ladder approximation uses a bare quark-gluon vertex, which violates QCD’s Slavnov-Taylor identity (STI). In order to overcome this deficiency, physicists are trying their best to go beyond the rainbow-ladder approximation. Much work has been done in this direction. Here we just name a few examples: the Ball-Chiu (BC) vertex derived from the vector Ward-Takahashi identity (WTI) \[18,19\], the CP vertex which takes into account some transversal effects \[20\] and the vertex derived from the transversal WTI \[21–23\], etc. As was shown in Ref. \[24\], if one deletes the ghost amplitudes and the gluon dressing function factor from the STI then the result has the form of a color matrix times the WTI structure. Here one notes that the BC vertex ansatz multiplied by the color matrix will satisfy such a relation. So in this paper we adopt such a quark-gluon vertex ansatz to explore the effect of vertex dressing on the tensor vacuum susceptibility.
When one tries to calculate the dressed quark propagator from the DSE using the BC vertex, one should construct a consistent kernel approximation corresponding to this vertex. How to construct systematic and convergent expansions for the kernels of DSE is a long-standing unsolved problem. Recently, an important progress in this problem has been achieved in Ref. [25]. The authors in Ref. [25] have proposed a Bethe-Salpeter kernel which is valid for a general quark-gluon vertex. This provides a theoretical foundation for calculating the tensor vacuum susceptibility beyond the rainbow-ladder approximation. In the present paper we shall use this method to calculate the tensor vacuum susceptibility.

In order to make this paper self-contained, let us first recall the definition of vacuum susceptibilities. In the QCD sum rule external field approach, the QCD vacuum susceptibility is tightly related to the linear response of the dressed quark propagator coupled nonperturbatively to an external current $J^\Gamma(y)\mathcal{V}_\Gamma(y) \equiv \bar{q}(y)\Gamma q(y)\mathcal{V}_\Gamma(y)$ ($q(y)$ is the quark field, $\Gamma$ stands for the appropriate combination of Dirac, flavor, color matrices and $\mathcal{V}_\Gamma(y)$ is the variable external field of interest) [1–3]. Here following Ref. [11], we adopt the following definition for the tensor vacuum susceptibility

$$\chi^Z = \left\{ \text{Tr}[\sigma_\eta \eta \zeta] \sigma \eta \zeta S\Gamma \cdot ZS - \text{Tr}[\sigma_\eta \eta \zeta S_0 \Gamma_0 \cdot S_0] \right\} / Z_{\eta \zeta}(0) : \bar{q}(0)q(0) : |0\rangle,$$

where $Z_{\eta \zeta}$ denotes the variable external tensor field, $\Gamma$ and $\Gamma_0 \ (\Gamma_0^{\mu \nu} = \sigma^{\mu \nu})$ denote the full and free tensor vertex, $S$ and $S_0$ are the full and free quark propagators. $\langle \bar{0} | : \bar{q}(0)q(0) : |0\rangle$ denotes the chiral quark condensate. Here it should be noted that Eq. (1) is essentially the tensor vacuum polarization, regularized by subtraction of the free vacuum polarization, and scaled by the scalar vacuum condensate. It can be obtained by external field differentiation of the propagator contracted with a bare vertex. Such a differentiation of a trace of a propagator, essentially a condensate, to produce a susceptibility as proportional to the associated vacuum polarization has recently been used by some of the same authors in Ref. [26] for the vector and axial-vector vacuum susceptibility, and also in Refs. [27, 28] for the scalar and pseudoscalar vacuum susceptibility.

From Eq. (1) we can see that the tensor vacuum susceptibility is closely related to the dressed quark propagator and the dressed tensor vertex at zero total momentum. Now we turn to the calculation of the dressed quark propagator and the dressed tensor vertex at zero total momentum in the DS approach. In the DS approach, the gap equation for the
dressed quark propagator $S$ in the chiral limit can be written as

$$S(p)^{-1} = Z_2 i \gamma \cdot p + \Sigma(p)$$

with

$$\Sigma(p) = Z_1 \int_q g^2 D_{\mu\nu}(p - q) \frac{\lambda^a}{2} \gamma_{\mu} S(q) \frac{\lambda^a}{2} \Gamma_{\nu}^0(q, p),$$

where $D_{\mu\nu}(k)$ is the dressed gluon propagator and $\Gamma_{\nu}(q, p)$ is the dressed quark-gluon vertex. The quark-gluon vertex and quark wave-function renormalization constants, $Z_{1,2}(\zeta^2, \Lambda^2)$, also depend on the gauge parameter.

The gap equation’s solution has the form

$$S(p)^{-1} = i \gamma \cdot p A(p^2, \zeta^2) + B(p^2, \zeta^2)$$

and the mass function $M(p^2) = B(p^2, \zeta^2)/A(p^2, \zeta^2)$ is renormalisation point independent. The quark propagator can be obtained from Eq. (2) with the following renormalisation condition (since QCD is asymptotically free, one can choose this renormalisation condition):

$$S(p)^{-1} \bigg|_{p^2 = \zeta^2} = i \gamma \cdot p.$$

The renormalized fully-dressed tensor vertex $\Gamma_{\mu\nu}$ satisfies an inhomogeneous Bethe-Salpeter equation:

$$\Gamma_{\mu\nu}(k, P; \zeta) = Z_T \sigma_{\mu\nu} + \int_q [S(q_+) \Gamma_{\mu\nu}(q, P) S(q_-)]_{sr} K_{ts}^{\mu\nu}(q, k; P).$$

Here $k$ is the relative and $P$ the total momentum of the quark-antiquark pair; $q_\pm = q \pm P/2$; $r, s, t, u$ represent colour and Dirac indices; and $K$ is referred to as the fully-amputated quark-antiquark scattering kernel. $Z_T$ is the renormalisation constant for the tensor vertex.

For the specific calculation of $\chi^Z$, one only requires the tensor vertex at $P = 0$. From general Lorentz structure analysis and the asymmetry of the tensor vertex $\Gamma_{\mu\nu}$ with respect to the indices $\mu$ and $\nu$, we can write down the general form of the tensor vertex

$$\Gamma_{\mu\nu}(p, 0) = \sigma_{\mu\nu} E(p^2) + (\gamma_\mu p_\nu - \gamma_\nu p_\mu) F(p^2) + i \gamma \cdot p (\gamma_\mu p_\nu - \gamma_\nu p_\mu) G(p^2).$$

Substituting Eqs. (4) and (7) into Eq. (1), we can obtain the final expression for calculating the tensor vacuum susceptibility

$$\chi^Z = \frac{3}{16\pi^2 a} \int_0^\infty dss \left\{ 2E(s) \left[ \frac{B(s)}{sA^2(s) + B^2(s)} \right]^2 - \frac{sG(s)}{sA^2(s) + B^2(s)} \right\},$$
where \( a = \langle \bar{q}(0) q(0) : \bar{q}(0) q(0) \rangle \) is the two-quark condensate. Here we note that in obtaining the above equation, we have made use of the fact that the subtraction term vanishes.

In phenomenological applications, one may proceed by considering the truncation scheme for the DSEs and BSEs, especially for the dressed gluon propagator, the dressed quark-gluon vertex and the four-point dressed quark-antiquark scattering kernel. The important information about the kernel of QCD’s gap equation can be phenomenologically drawn by a dialogue between DSE studies and results from numerical simulations of lattice-regularized QCD \([29–32]\). The ansatz that is typically implemented in the quark propagator’s gap equation can be written as

\[
Z_1 g^2 D_{\rho\sigma}(p-q) \Gamma_{\sigma}(q,p) \to \mathcal{G}((p-q)^2) D^\text{free}_{\rho\sigma}(p-q) \frac{\alpha}{2} \Gamma_{\sigma}(q,p),
\]

wherein \( D^\text{free}_{\rho\sigma}(\ell) \) is the Landau-gauge free gauge-boson propagator, \( \mathcal{G}(\ell^2) \) is a model effective-interaction and \( \Gamma_{\sigma}(q,p) \) is a vertex ansatz.

Over the past few years, the most usually used approximation is the rainbow-ladder approximation \([12–17]\), where the dressed quark-gluon vertex \( \Gamma_{\mu}(q,p) \) is replaced by the bare vertex \( \gamma_{\mu} \), and in the BS equation the ladder kernel is used. Rainbow-ladder approximation is the lowest order truncation scheme for the DSE. It is the nonperturbative symmetry-conserving truncation scheme because it satisfies the axial-vector WTI. Models formulated using the rainbow-ladder DSE to describe the quark dynamics within hadrons were found to provide good and compact descriptions of the light pseudoscalar and vector mesons. However, the rainbow-ladder DSE cannot describe well the properties of scalar mesons. So physicists are trying to go beyond the rainbow-ladder approximation for years. The key points to go beyond the rainbow-ladder approximation are the dressed quark-gluon vertex and the four-point quark-antiquark scattering kernel.

For the dressed quark-gluon vertex, we can employ the BC vertex \([18–19]\)

\[
i \Gamma_{\sigma}(k, \ell) = i \Sigma_A(k^2, \ell^2) \gamma_{\sigma} + (k + \ell)_{\sigma} \times \left[ \frac{i}{2} \gamma \cdot (k + \ell) \Delta_A(k^2, \ell^2) + \Delta_B(k^2, \ell^2) \right],
\]

where

\[
\Sigma_F(k^2, \ell^2) = \frac{1}{2} \left[ F(k^2) + F(\ell^2) \right], \quad \Delta_F(k^2, \ell^2) = \frac{F(k^2) - F(\ell^2)}{k^2 - \ell^2},
\]

with \( F = A, B \), viz., the scalar functions in Eq. \((4)\). Here it should be noted that the BC vertex satisfies the vector WTI.
Now one should find a kernel consistent with the BC vertex ansatz. This is a difficult task that many scientists try to do. Recently, great progress has been done on this aspect. The authors in Ref. [25] have found a way to constrain the kernel for the general vertex. Following their method, an exact form of the inhomogeneous BSE for the tensor vertex $\Gamma_{\mu\nu}(k,0)$ can also be written as

$$
\Gamma_{\mu\nu}(k,0) = Z_T\sigma_{\mu\nu} - \int_q g^2 D_{\alpha\beta}(k-q) \frac{\lambda^a}{2} \gamma_{\alpha} S(q) \Gamma_{\mu\nu}(q,0) S(q) \frac{\lambda^a}{2} \Gamma(\beta, k) + \int_q g^2 D_{\alpha\beta}(k-q) \frac{\lambda^a}{2} \gamma_{\alpha} S(q) \frac{\lambda^a}{2} \Lambda_{\mu\nu\beta}(k, q; 0),
$$

(12)

where $\Lambda_{\mu\nu\beta}(k, q; 0)$ is a four-point Schwinger function that is completely defined via the quark self-energy [33, 34]. It satisfies the similar identity as those in Ref. [25]

$$(k - q)_{\beta i} \Lambda_{\mu\nu\beta}(k, q; 0) = \Gamma_{\mu\nu}(k, 0) - \Gamma_{\mu\nu}(q, 0).$$

(13)

Then we can obtain

$$
i \Lambda_{\mu\nu\beta}(k, q; 0) = 2l_{\beta}[\Delta E(q, k; 0) + (\gamma_\mu l_\nu - \gamma_\nu l_\mu) \Delta_F(q, k; 0)]$$

$$+ (\gamma_\mu \delta_\nu\beta - \gamma_\nu \delta_\mu\beta) \Sigma_F(q, k; 0) + 2l_\gamma \cdot l(\gamma_\mu l_\nu - \gamma_\nu l_\mu) \Delta_F(q, k; 0)$$

$$+ (\gamma \cdot l(\gamma_\mu \delta_\nu\beta - \gamma_\nu \delta_\mu\beta) \Sigma_F(q, k; 0) + \gamma_\beta(\gamma_\mu l_\nu - \gamma_\nu l_\mu) \Sigma_G(q, k; 0)$$

$$+ \frac{1}{4} \gamma_\beta(k^2 - q^2)[\gamma_\mu(q - k)_\nu - \gamma_\nu(q - k)_\mu] \Delta_G(q, k; 0).$$

(14)

Herein we employ a simplified form of the renormalisation-group-improved effective interaction proposed in Refs. [12–17]; viz., we retain only that piece which expresses the long-range behavior ($s = k^2)$:

$$
\frac{\mathcal{G}(s)}{s} = \frac{4\pi^2}{\omega^6} D s e^{-s/\omega^2}.
$$

(15)

This is a finite width representation of the form introduced in Ref. [36], which has been rendered as an integrable regularisation of $1/k^4$ [37]. Equation (15) delivers an ultraviolet finite model gap equation. Hence, the regularisation mass-scale can be removed to infinity and the renormalisation constants set equal to one.

The active parameters in Eq. (15) are $D$ and $\omega$ but they are not independent. In reconsidering a renormalisation-group-improved rainbow-ladder fit to a selection of ground state observables [13], Ref. [15] noted that a change in $D$ is compensated by an alteration of $\omega$. This feature has further been elucidated and exploited in Refs. [16, 17, 35]. For
FIG. 1: Dressed quark propagator. Left panel – $A(p^2)$, right panel – $B(p^2)$. In both panels, Dashed curve: calculated in rainbow-ladder truncation; solid curve: calculated with BC vertex ansatz.

$\omega \in [0.3, 0.5]\text{ GeV}$, with the interaction specified by Eqs. (9), (10) and (15), fitted in-vacuum low-energy observables are approximately constant along the trajectory

$$\omega D = (0.8\text{ GeV})^3 =: m_3^3.$$  

Herein, we employ $\omega = 0.5\text{ GeV}, D = m_3^3/\omega = 1.0 \text{ GeV}^2$.

So now with the BC vertex ansatz and the model effective interaction, the equations of the DSE for the dressed quark propagator and the BSE for the dressed tensor vertex are reduced to a closed system of equations. We can numerically calculate them with iteration method. In Fig. 1 we plot the functions obtained through solving the gap equation and in Fig. 2 those which describe the dressed tensor vertex.

It is apparent in Fig. 1 that the vertex Ansatz has a quantitative impact on the magnitude and point-wise evolution of the gap equation’s solution. That this should be anticipated is plain from Ref. [38]. Moreover, the pattern of behavior can be understood from Ref. [39]: the feedback arising through the $\Delta_B$ term in the BC vertex, Eq. (10), absent in the rainbow approximation, always acts to alter the domain upon which $A(p^2)$ and $M(p^2)$ differ significantly in magnitude from their respective free-particle values. Since $E(p^2)$, $F(p^2)$ and $G(p^2)$ are derived quantities, their behavior does not require explanation. We plot the integrand in Eq. (8) in Fig. 3 for each vertex ansatz. From the figure we can see that there is no far-ultraviolet tail in the integrand so that the we do not need regularization here. The
resulting tensor vacuum susceptibilities are

\[ \chi_{BC}^Z = 0.05573 \text{ GeV}^{-1}, \quad \chi_{RL}^Z = 0.08672 \text{ GeV}^{-1}. \]  

(17)

The above result shows that the numerical value of the tensor vacuum susceptibility obtained in the BC vertex approximation is much smaller than that in the rainbow-ladder approximation. Here it should be noted that in the above calculations of tensor vacuum susceptibility using the effective interaction (15) in the rainbow-ladder truncation and the BC vertex, we have chosen the same model parameters for the effective interaction. As is shown in Ref. [25], the amount of chiral symmetry breaking (as measured by the chiral condensate) and related quantities such as the pion decay constant are very different between these two truncation schemes. Therefore, when calculating the tensor vacuum susceptibility employing the BC vertex, a reasonable approach is to use refitted model parameters in the effective interaction (15) in the calculation. Because the active parameters \( D \) and \( \omega \) in Eq. (15) are not independent, one can refit the model parameters from one physical quan-
tity, for example, the chiral condensate. Under the BC vertex, the value of the parameter $D$ fitted from the chiral condensate is $D = \frac{1}{2} \text{GeV}^2$ (see Ref. [27]). The results for the dressed quark propagator, the scalar functions $E(p^2)$, $F(p^2)$, $G(p^2)$, and the integrand in Eq. (14), calculated from both the rainbow-ladder truncation and the BC vertex with refitted model parameters are shown in Figs. 4 to 6. With refitted parameters in the BC vertex approximation, the resulting tensor vacuum susceptibility are

$$\chi^{Z}_{BC} = 0.07886 \text{GeV}^{-1}$$

$$\chi^{Z}_{RL} = 0.08672 \text{GeV}^{-1}.$$  

(18)

So, compared with the rainbow-ladder truncation result, the value of $\chi^{Z}$ in the BC vertex approximation is reduced by about 10%. Therefore, one can draw the conclusion that in the calculation of the tensor vacuum susceptibility in the framework of the DS approach the dressing effect of the quark-gluon vertex is not large.

In Fig. 7 we depict the evolution of the tensor vacuum susceptibility with increasing interaction strength, $\mathcal{I} = D/\omega^2$. The behavior may readily be understood. For $\mathcal{I} = 0$ one has a noninteracting theory and the “vacuum” is unperturbed by the external tensor field. Hence, the susceptibility is zero. The tensor vacuum susceptibility remains zero until the interaction strength $\mathcal{I}$ reaches a critical value, $\mathcal{I} = \mathcal{I}_c$. When $\mathcal{I} > \mathcal{I}_c$, the tensor vacuum susceptibility becomes larger quickly and then goes down slowly for both the rainbow-ladder approximation and the BC vertex approximation. Those critical values for the interaction strength are: $\mathcal{I}_c^{RL} = 1.93$, $\mathcal{I}_c^{BC} = 1.41$. It can be seen that the critical point in the rainbow-ladder approximation is larger than that in the BC vertex approximation. This is easy to understand, because the effect of the BC vertex itself amounts to enhancing the interaction
FIG. 4: Dressed quark propagator. *Left panel* – $A(p^2)$, *right panel* – $B(p^2)$. In both panels, *dashed curve*: calculated in rainbow-ladder truncation; *solid curve*: calculated with BC vertex ansatz with refitted model parameters.

FIG. 5: $P = 0$ scalar vertex, Eq. (7): *upper left panel* – $E(p^2)$, *upper right panel* – $F(p^2)$, *lower panel* – $G(p^2)$. In all panels, *dashed curve*: calculated in rainbow-ladder truncation; *solid curve*: calculated with BC vertex ansatz with refitted model parameters.
FIG. 6: Integrand in Eq. (8) – Dashed curve: calculated in rainbow-ladder truncation; solid curve: calculated with BC vertex ansatz with refitted model parameters.

FIG. 7: Dependence of the chiral susceptibility on the interaction strength in Eq. (15); viz., $I := D/\omega^2$: dashed curve, RL vertex; solid curve, BC vertex.

strength. The authors in Ref. [27] has explained the nature of the critical interaction strength which denotes a second-order phase transition.

For $I < I_c$, the interaction strength is not sufficient to generate a non-zero scalar term in the dressed quark self-energy in the chiral limit. That means below the critical value, dynamical chiral symmetry breaking is impossible. The situation changes at $I_c$, for $I > I_c$ a $B \neq 0$ solution is always possible. Moreover, when $I < I_c$, the interaction strength is also not sufficient to generate the non-zero $F$ and $G$ functions in the dressed tensor vertex. That is the reason why the tensor vacuum susceptibility remains zero when $I < I_c$.

To summarize, using the expression obtained in the QCD sum rule external field approach in Ref. [11], we extend the calculation of tensor vacuum susceptibility in the rainbow-ladder approximation of the DS approach in Ref. [11] to that of employing the BC vertex
approximation. Here a key problem is how to construct a consistent Bethe-Salpeter kernel for a dressed quark-gluon vertex ansatz whose diagrammatic content is unknown. Recently, significant progress in this problem was achieved in Ref. [25]. In this paper, following the work of Ref. [25], we construct the kernel for the dressed tensor vertex at $P = 0$ which is needed in the calculation of tensor vacuum susceptibility. Then we perform a consistent calculation of the tensor vacuum susceptibility beyond the rainbow-ladder approximation. Our results show that compared with its rainbow-ladder approximation value, the tensor vacuum susceptibility in the BC vertex approximation is reduced by about 10%. This shows that the dressing effect of the quark-gluon vertex is not large in the calculation of the tensor vacuum susceptibility in the framework of the DS approach. In this paper we also demonstrate that the tensor vacuum susceptibility can be used to demarcate the domain of coupling strength within a theory upon which chiral symmetry is dynamically broken. For couplings below the associated critical value and in the absence of confinement, the tensor vacuum susceptibility remains zero. This situation changes until the interaction strength is larger than a critical point. It is found that the critical point in the rainbow-ladder approximation is larger than that in the BC vertex approximation. This is easy to understand, because the effect of the BC vertex itself amounts to enhancing the interaction strength.

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