On the massless contributions to the vacuum polarization of heavy quarks

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Abstract

In Ref. [1] Groote and Pivovarov have given notice of a possible fault in the use of sum rules involving two-point correlation functions to extract information on heavy quark parameters, due to the presence of massless contributions that invalidate the construction of moments of the spectral densities. Here we show how to circumvent this problem through a new definition of the moments, providing an infrared safe and consistent procedure.

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1 Introduction

The vacuum polarization function suitable for extracting fundamental information of heavy quark-antiquark systems is built from the electromagnetic current \( j^\mu(x) = e_Q \bar{Q}(x) \gamma^\mu Q(x) \) of the heavy quark \( Q \) of mass \( M \) and \( e_Q \) electric charge:

\[
\Pi_{\mu\nu}(q) = i \int d^4x \, e^{iqx} \langle 0 | T j_\mu(x) j^\nu_v(0) | 0 \rangle = \left( -g_{\mu\nu} q^2 + q_\mu q_\nu \right) \Pi(q^2). \tag{1}
\]

As it is well known two–point functions are analytic except for singularities at simple poles or branch cuts, the latter being originated by normal thresholds of production of internal on–shell states. Implicitly assuming that the absorptive part of \( \Pi(q^2) \) starts at the massive two–particle threshold \( q^2 = 4M^2 \), vanishing below this point, the correlator satisfies the once–subtracted dispersion relation \[2\]:

\[
\hat{\Pi}(q^2) = \Pi(q^2) - \Pi(0) = \frac{q^2}{\pi} \int_{4M^2}^\infty \frac{ds}{s} \frac{\text{Im} \, \Pi(s)}{s - q^2 - i\epsilon}. \tag{2}
\]

This dispersion relation has been extensively used to determine heavy quark parameters within the method of sum rules because it allows to relate experimental input, on the right–hand side, with theoretical perturbative evaluations on the left–hand side \[3\]. Indeed \( \text{Im} \, \Pi(q^2) \) refers to the total cross section of heavy quark production \( \sigma(e^+e^- \rightarrow QQ) \). In practice, the spectral density \( \text{Im} \, \Pi(s) \) is poorly known experimentally at very high energies and, in addition, we are interested in the very low energy domain because it is more sensitive to the heavy quark mass. Therefore one uses derivatives of the vacuum polarization at the origin, called moments, to be responsive to the threshold region:

\[
\mathcal{M}_n = \frac{1}{n!} \left( \frac{d}{dq^2} \right)^n \Pi(q^2)|_{q^2=0}. \tag{3}
\]

Until present the evaluation of the perturbative two–point correlation function \( \Pi(q^2) \) has only been carried out completely, with massive quarks, up to \( \mathcal{O}(\alpha_s^2) \) \[4\] and the procedure above has been termed consistent and effective in its task because the first branch point is set at the massive two–particle threshold. However in Ref. \[1\] Groote and Pivovarov have pointed out that at \( \mathcal{O}(\alpha_s^3) \) there is a contribution to the correlator which contains a three–gluon massless intermediate state (see Fig. 1(a)). Its absorptive part starts at zero energy and, therefore, Eq. (2) is no longer correct. Moreover those authors have also warned about the fact that, at this perturbative order, the massless intermediate state invalidates the definition of the moments \( \mathcal{M}_n \) for \( n \geq 4 \) because they become singular. In Ref. \[5\] an infrared safe redefinition of the moments, to cure the latter problem, has been provided; it consists in evaluating the moments at an Euclidean point \( q^2 = -s_E, s_E > 0 \), thus avoiding the singular behaviour. Nevertheless the fault in Eq. (2) due to the massless threshold still represents a problem because even if, as we will justify later on, we substitute this dispersion relation by

\[
\hat{\Pi}(q^2) = \frac{q^2}{\pi} \int_{4M^2}^\infty \frac{ds}{s} \frac{\text{Im} \, \Pi_{QQ}(s)}{s - q^2 - i\epsilon} + \frac{q^2}{\pi} \int_0^\infty \frac{ds}{s} \frac{\text{Im} \, \Pi_{3g}(s)}{s - q^2 - i\epsilon}, \tag{4}
\]

\[1\] Sometimes de Adler function defined as \( \partial \Pi(q^2)/\partial \ln q^2 \), to get rid of the subtraction constant, is used.
Figure 1: (a) $O(\alpha_s^3)$ diagram contributing to the vacuum polarization function of the heavy quark current (the vertical dashed line indicates the massless cut). (b) “Effective” diagram obtained by integrating out the fermion loops. It also has the topological structure of the “reduced” diagram that determines the massless cut singularity.

(where the notation is self-explicative), the spectral function $\text{Im} \Pi_{3g}(s)$ associated to the cut in Fig. 1(a) would hardly be implemented phenomenologically as gluons hadronize to both heavy and light quark pairs. Perturbatively the three–gluon cut would contribute to $Q\bar{Q}$ production, i.e. to $\text{Im} \Pi_{Q\bar{Q}}(q^2)$, but at higher order in $\alpha_s$. Therefore if we attach to an $O(\alpha_s^3)$ sum rule analysis, that contribution should be extracted from the perturbative $\Pi(q^2)$ evaluation. In this note we provide a bypass to recover the balance between the right-hand and left-hand parts of Eq. (4).

2 Moments and the massless cut

The perturbative contribution given by the diagram in Fig. 1(a) has been calculated at small $q^2$ ($q^2 \ll M^2$) in Ref. [1]. In this limit the quark triangle loop can be integrated out and it ends up in the diagram in Fig. 1(b) generated by an induced effective current describing the interaction of the vector current with three gluons $\pi^2$.

$$ J^\mu = -\frac{\pi}{180M^4} \left( \frac{\alpha_s}{\pi} \right)^2 (5 \partial_\nu \mathcal{O}_1^{\mu\nu} + 14 \partial_\nu \mathcal{O}_2^{\mu\nu}) , $$

with

$$ \mathcal{O}_1^{\mu\nu} = d_{abc} G_a^{\mu\nu} G_b^{\alpha\beta} G_c^{\alpha\beta} , $$

$$ \mathcal{O}_2^{\mu\nu} = d_{abc} G_a^{\mu\alpha} G_b^{\beta\beta} G_c^{\beta\nu} , $$

where $G_a^{\mu\nu}$ is the gluon strength field tensor. The effective current in the QED case ($G_a^{\mu\nu} \to F^{\mu\nu}, \alpha_s \to \alpha_{em}, d_{abc} \to 1$) can be easily identified from the lowest order Euler-Heisenberg Lagrangian (see Ref. [3]).

2 The permutations of the three gluons in Fig. 1(a) are already included in the definition of the effective current.
The correlator of the induced current (5) is then evaluated in the configuration space giving:

$$\langle 0 | T J_\mu(x) J^\dagger_\nu(0) | 0 \rangle = -\frac{34}{2025 \pi^3 M^8} (\frac{\alpha_s}{\pi})^3 d_{abc}d_{abc} \left( \partial_\mu \partial_\nu - g_\mu\nu \partial^2 \right) \frac{1}{x^{12}}.$$  \hspace{1cm} (7)

In momentum space we need to perform the Fourier transform of Eq. (7). Following the differential regularization procedure [6], which works directly in configuration space, the result for the vacuum polarization contribution of the diagram in Fig. 1(b) at small $q^2$ reads

$$\Pi_{\mu\nu}(q) = \frac{17}{2916000\pi^2} d_{abc}d_{abc} \left( \frac{\alpha_s}{\pi} \right)^3 (q_\mu q_\nu - q^2 g_\mu\nu) \left( \frac{q^2}{4M^2} \right)^4 \ln \left( \frac{\mu^2}{-q^2} \right) + \mathcal{O}\left[ (\frac{q^2}{M^2})^5 \right],$$  \hspace{1cm} (8)

with $\mu$ the renormalization point in this scheme, and $d_{abc}d_{abc} = 40/3$.

As noticed by Groote and Pivovarov [1], moments associated to the diagram in Fig. 1(b) are not defined if $n \geq 4$. Indeed differentiating Eq. (8) four times, at $q^2 \approx 0$, we get:

$$\frac{1}{4!} \left( \frac{d}{dq^2} \right)^4 \Pi(q^2)|_{q^2=0} = \frac{17}{2187000\pi^2} \left( \frac{\alpha_s}{\pi} \right)^3 \left( \frac{1}{4M^2} \right)^4 \left[ \ln \left( \frac{\mu^2}{-q^2} \right) - \frac{25}{12} \right] + \mathcal{O}\left[ \frac{q^2}{M^{10}} \right],$$  \hspace{1cm} (9)

whose real part clearly diverges if we set $q^2 = 0$. Larger $n$ moments are also infrared divergent, and so the authors of Ref. [1] conclude that the standard sum rule analysis must limit the accuracy of theoretical calculations for the $n \geq 4$ moments to the $\mathcal{O}(\alpha_s^2)$ order of perturbation theory.

One obvious way out of this infrared problem is to avoid the $q^2 = 0$ point. As it has been discussed in Ref. [3], this solution is rather ill-conditioned from the phenomenological side though. Moreover we notice (as commented earlier) that it is not possible to implement, from the available experimental information, the second term in the right-hand side of Eq. (4). However, if one does not insist in using full vacuum polarization for the sum rule analysis there is a way to overcome this infrared problem.

### 3 Infrared safe definition of the moments

The study of analytic properties of perturbation theory amplitudes shows that their singularities are isolated and, therefore, we can discuss each singularity of a perturbative amplitude by itself [3]. As a consequence, $\Pi(q^2)$ in Eq. (1) satisfies a dispersion relation from Cauchy’s theorem [3]:

$$\Pi(q^2) = \sum_n \frac{q^2}{2\pi i} \int_{s_n}^\infty ds \frac{[\Pi(s)]_n}{s - q^2 - i\epsilon}. \hspace{1cm} (10)$$

Here $[\Pi(s)]_n$ provides the discontinuity across a branch cut starting at the branch point $s_n$.

\[3\]This expression also gives the residue $R_i$ of a pole at $s = s_i$ if we interpret the discontinuity as $\text{Im}\Pi_i = \pi R_i \delta(s - s_i)$. However we do not consider the existence of $Q\bar{Q}$ Coulomb bound states, as it is not relevant for our discussion.
In the perturbative calculation, every discontinuity function $[\Pi(s)]_n$ can be associated to a “reduced” Feynman diagram obtained by contracting internal off–shell propagators to a point and leaving internal on–shell lines untouched. Its contribution is written down following the Cutkosky rules for the graph. The reduced diagram corresponding to the massless cut in Fig. 1(a) has the topological structure of the part (b) of that Figure. Let us emphasize though that our following discussion is not grounded on the $q^2 \ll M^2$ regime where the fermion loops have been integrated out: the reduced diagram is just a symbol that specifies a singularity, and the black dots in Fig. 1(b) keep all the analytical structure of the fermion loops.

In a general diagram the discontinuity across a specified cut needs not to be a pure real function in the physical region, only the sum of all cuts in a given channel gives the total imaginary part. Hence the separation between the imaginary parts coming from different final states, as performed in Eq. (4) for the vacuum polarization, does not seem to come directly from the Cutkosky rules. Nevertheless in the heavy quark correlator the discontinuity across the three–gluon cut gives a contribution to the spectral function that is unequivocally real:

$$\frac{1}{2i} [\Pi(s)]_{3g} = \text{Im} \Pi_{3g}(s) = -\frac{1}{6s} \int dR_{3g} \langle 0 | j^\mu | 3g \rangle \langle 3g | j^\mu_\dagger | 0 \rangle ,$$

from which the dispersive part can be evaluated independently of the $Q\overline{Q}$ cuts\footnote{The integration in Eq. (11) extends to the available three–gluon phase space.}. Accordingly we conclude that we can identify and isolate the troublesome massless cut contribution to the two–point function. Indeed Eqs. (11) and (12) justify our previous Eq. (4). This assertion might seem obvious but it is not: A $Q\overline{Q}$ cut on the right–hand fermion loop in Fig. 1(a) does not provide, by itself, a pure real contribution. Only when both $Q\overline{Q}$ cuts, on the left–hand and right–hand fermion loops of Fig. 1(a), are added we get $\text{Im} \Pi_{Q\overline{Q}}$.

Let us go back then to Eq. (4). All the difficulty with the phenomenological application of the sum rules is now the fact that the contribution from the three–gluon cut is contained in both sides of the equality. Thus we propose an infrared safe definition of the moments by the trivial subtraction:

$$\hat{\Pi}_{Q\overline{Q}}(q^2) \doteq \hat{\Pi}(q^2) - \frac{q^2}{\pi} \int_0^\infty ds \frac{\text{Im} \Pi_{3g}(s)}{s - q^2 - i\epsilon} = \frac{q^2}{\pi} \int_{4M^2}^\infty ds \frac{\text{Im} \Pi_{Q\overline{Q}}(s)}{s - q^2 - i\epsilon} ,$$

$$\hat{\mathcal{M}}_n \doteq \mathcal{M}_n - \frac{1}{\pi} \int_0^\infty ds \frac{\text{Im} \Pi_{3g}(s)}{s^{n+1}} .$$

Of course Eqs. (12) and (13) are meaningless unless we give a precise prescription about how to subtract the contribution of the massless cuts represented by $\text{Im} \Pi_{3g}$. Our previous discussion gives us the tool to proceed. Once the full $\mathcal{O}(\alpha_s^3) \Pi(s)$ is calculated we can extract the imaginary part starting at $s = 0$ (which should go with a $\theta(s)$ function) for any value of $s$. It is clear that the $\theta(s)$ and $\theta(s - 4M^2)$ terms in the imaginary part of the vacuum polarization function correspond to three–gluon massless and to $Q\overline{Q}$ cut graphs, respectively, and $\text{Im} \Pi_{3g}$ and $\text{Im} \Pi_{Q\overline{Q}}$ are easy to distinguish, as Eq. (11) prevents the appearance of mixed $\theta(s) \cdot \theta(s - 4M^2)$ terms. Therefore we identify $\text{Im} \Pi_{3g}$ and we now plug it in the dispersion integral of the right–hand
Figure 2: Examples of perturbative non-heavy quark current correlators at $O(\alpha_s^2)$ (a) and $O(\alpha_s^3)$ (b) that contribute to the production of $Q\bar{Q}$ states.

side of Eq. (13) and perform such integration. Divergences contained in both this integral and $M_n$ as $q^2 \to 0$ will cancel with each other if the same infrared regularization is employed in the two quantities. The intuitive choice would be a low-energy cutoff $s_0 > 0$, and Eq. (13) would be more precisely written as:

$$\tilde{M}_n \equiv \lim_{s_0 \to 0^+} \left[ \frac{1}{n!} \left( \frac{d}{dq^2} \right)^n \Pi(q^2)|_{q^2 = -s_0} - \frac{1}{\pi} \int_0^\infty ds \frac{\text{Im} \Pi_{3g}(s)}{(s + s_0)^n} \right],$$

(14)

where a vanishing term in the $s_0 \to 0^+$ limit has been omitted.

The evaluation of the $M_n$ moments at $q^2 = 0 < 4M^2$ made sense because, up to $O(\alpha_s^2)$, this point is unphysical and the moments are well defined through an analytic continuation from the high–energy region. However note that the absorptive three–gluon contribution starts at $q^2 = 0$ where perturbative QCD is unreliable. This introduces a further new difficulty in evaluating $M_n$ moments at $q^2 = 0$, as we reach the physical non–perturbative region. Our definition of the moments, $\tilde{M}_n$, in Eq. (13), skips this problem by fully eliminating the massless terms and, therefore, the final heavy quark sum rule will only involve physics at $q^2 > 4M^2$.

The method discussed in this Section could be extended to more general two–point correlators involving intermediate $Q\bar{Q}$ states in their perturbative expansion. Indeed the correlator of two light quark vector currents has $O(\alpha_s^2)$ contributions with an internal loop of heavy quarks (Fig. 3 (a)) and, similarly, the asymmetric correlator of a heavy and a light vector quark currents is no longer vanishing at $O(\alpha_s^3)$ (Fig. 3 (b)). The absorptive part coming from the $Q\bar{Q}$ cuts in the previous examples contribute to the phenomenological input $\sigma(e^+e^- \to Q\bar{Q})$ in the usual sum rule analysis. Correspondingly they should be accounted for in the theoretical side. In short, the production of $Q\bar{Q}$ states concerns not only the correlator of a couple of heavy quark currents and, for a more rigorous use of the sum rules method, this imbalance should be taken into account and properly fixed. A two–point function built from the sum of the electromagnetic currents associated to each quark flavour could be used in a generalized
version of Eq. (1): 

$$\Pi_{\mu\nu}^G(p) = i \int d^4x e^{ipx} \sum_{q,q'} e_q e_{q'} \langle 0\mid T(\bar{q}(x)\gamma_{\mu}q(x)) (\bar{q}'(0)\gamma_{\nu}q'(0)) \mid 0 \rangle ,$$  

(15)

where now $q$ and $q'$ stand for heavy or light quarks indistinctly, with electric charges $e_q$ and $e_{q'}$, and $q = q'$ is also allowed. As the different absorptive cuts contribute additively to Im$\Pi(p^2)$, the unwanted light quark and gluon $q\bar{q}$, $q\bar{q}g$, $ggg$, ... cuts could be identified and, through the dispersive technique, subtracted from the full $\Pi(p^2)$ result. Consequently we are left with every possible $Q\bar{Q}$ intermediate state arising from vector current production. Notwithstanding, the feasibility of this procedure from a technical point of view appears rather cumbersome and, at present, the experimental accuracy in the measurement of $\sigma(e^+e^- \rightarrow Q\bar{Q})$ cannot accommodate the corrections just discussed.

4 Discussion on QCD sum rules applications

The general rule given above is valid for all orders of perturbation theory, but it strongly relies in our ability to extract the massless absorptive part from the full result of $\Pi(q^2)$ calculated at a definite order. Beyond $\mathcal{O}(\alpha_s^2)$ complete analytical results for the heavy quark correlator would be cumbersome and only numerical approaches may be at hand. In this sense, it would be convenient to have a method to calculate Im$\Pi_{Q\bar{Q}}$ only based on Feynman graphs. We have already sketched such a method in the discussion following Eq. (10): we just need to sum up all the massless cut graphs to get Im$\Pi_{3g}$, and then proceed with the dispersion integration that gives the associated dispersive part [8]. For example, at $\mathcal{O}(\alpha_s^3)$, the only massless absorptive part comes from the three–gluon cut in the diagram of Fig. 1(a); let us call $M_{3g}^\mu$ the amplitude producing three gluons from the heavy quark current at lowest order (i.e. through the quark triangle loop in Fig. 3). The massless contribution to the absorptive part of the correlator is then:

$$\text{Im} \Pi_{3g}(s) = -\frac{1}{6s} \int dR_{3g} M_{3g}^\mu \cdot M_{3g}^* ,$$  

(16)

with the three–gluon phase space integral defined as

$$\int dR_{3g} \equiv \frac{1}{3!} \frac{\pi^2}{(2\pi)^5} \frac{4s}{s_1} \int_0^s ds_1 \int_0^{s-s_1} ds_2 ,$$  

(17)

in terms of the invariants $s_1 \equiv (k_1 + k_2)^2 = (q - k_3)^2$ and $s_2 \equiv (k_2 + k_3)^2 = (q - k_1)^2$, and $k_i$ being the momenta of the gluons. The real part would be obtained by integrating Eq. (16):

$$\frac{s_0}{\pi} \int_0^\infty dS \frac{\text{Im} \Pi_{3g}(s)}{s + s_0} = \frac{-s_0}{288(2\pi)^4} \int_0^\infty ds \frac{1}{s^3(s + s_0)} \int_0^s ds_1 \int_0^{s-s_1} ds_2 M_{3g}^\mu \cdot M_{3g}^* ,$$  

(18)

which, in principle, could be performed also numerically. The nth-derivative of relation (18) respect to $s_0$, in the limit $s_0 \rightarrow 0^+$, would give the infrared divergent contribution that should be subtracted from the full moments, as dictated by Eq. (14).
This evaluation solves the extraction procedure on the theoretical side. However the use of moments on heavy quark sum rules involves global quark–hadron duality which translates into the supposed equivalence between the theoretical evaluation of the moments and the phenomenological input. The implementation of the latter is not a trivial task. The total experimental cross section $\sigma(e^+e^- \to hadrons)$ can be split into two disjoint quantities: the cross section for producing hadrons with Q–flavoured states, and the production of hadrons with no Q–flavoured components. If the experimental set up was accurate enough to classify events into one of these two clusters, the first class would be the required ingredient for the phenomenological part of the heavy quark sum rule. However this separation, implemented in the theoretical side within perturbative QCD, is rather involved. Up to $\mathcal{O}(\alpha_s^2)$ there has not been any doubt, in the literature, that contributions to this side arise wholly from $Q\bar{Q}$ cuts in the heavy quark correlator $\Pi_{Q\bar{Q}}$. The physical picture behind this assertion relies in the assumption of factorization between hard and soft regions in the quark production process and subsequent hadronization. The hard region described with perturbative QCD entails the production of the pair of heavy quarks, and the soft part of the interaction is responsible for the observed final hadron content. Although possible, annihilation of the partonic state $Q\bar{Q}$ due to the later interaction is very unlikely, as jets arising from the short distance interaction fly apart before long–distance effects become essential. Consequently, each jet hadronizes to a content of Q–flavoured states with unit probability. As local duality is implicitly invoked, this picture is assumed to hold at sufficient high energies; hence perturbative corrections to the hard part are successively included through the heavy quark currents correlator.

The same physical picture applied to the three-gluon cut does not allow us to conclude whether this piece should be included or not in the theoretical side of the sum rules. Confinement tell us that this intermediate state hadronizes completely into hadrons with a content of light and/or heavy quarks indistinctly. It is conspicuous that if we could disentangle the heavy quark hadronization, $3g \to Q\bar{Q}$, we should include only this piece into the sum rule. Then the singularity at $q^2 = 0$ would disappear because heavy quarks are produced starting at $q^2 = 4M^2$. However there is no way to sort out light and heavy quark production off three gluons and,
therefore, if we extract this contribution from the heavy quark sum rules we are introducing an
incertitude in the procedure because we make sure that there is no light quark hadronization
but we miss the heavy quark production. It is easy to see that the induced error is small,
due to the fact that three gluons hadronize mostly to light hadrons. On one side, in the very
high energy region and following perturbative QCD with $N_F = 4$, we have only a $1/4 = 25\%$
probability of finding a specified pair of heavy quarks produced. And this is a generous upper
limit because when we go down in energy, phase space restrictions severely reduce the counting
of heavy quarks. Hence we estimate that excluding the three–gluon cut we introduce a tiny
very few percent error in the sum rules procedure.

5 Conclusions

We have shown that rigorous and straightforward results of the general theory of singu-
larities of perturbation theory amplitudes provide all–important tools to extract the unwanted
$\mathcal{O}(\alpha_s^3)$ three–gluon massless cut pointed out by Groote and Pivovarov from the vector current
correlator of heavy quarks. We conclude that the appropriate procedure to obtain information
about the heavy quark parameters should make use of the infrared safe corrected moments,
deﬁned in Eq. (14), that now indeed satisfy the modiﬁed sum rule:

$$
\tilde{M}_n = \frac{1}{\pi} \int_{4M^2}^{\infty} ds \frac{\text{Im}\Pi_{qq}(s)}{s^{n+1}},
$$

where the right–hand side can be extracted from the heavy quark production cross section
$\sigma(e^+e^- \rightarrow Q\bar{Q})$.

Finally we have pointed out that, starting already at $\mathcal{O}(\alpha_s^2)$, the use of sum rules associated
to heavy vector current correlators shows and imbalance between the phenomenological input
in the dispersion relation and the perturbative two–point function. This is due to the fact that
the cross section of production of $Q\bar{Q}$ heavy quarks is contained not only in the correlator of
two heavy quark currents but in those involving light quarks too. We have indicated how to
improve the application of sum rules by constructing a generalized correlator of vector currents
in Eq. (15) and, afterwards, extracting all the perturbative information not related with the
production of heavy quarks, as we did in detail for the three–gluon cut.

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