Constraints on the finite-density spectral densities 
of vector channel

Xuemin Jin

Department of Physics and Center for Theoretical Physics
University of Maryland, College Park, Maryland 20742

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Abstract

Sum rules for the variation of finite-density spectral density of vector channel with baryon density are derived based on dispersion relations and the operator product expansion. These sum rules may serve as constraints on the phenomenological models for the finite-density spectral densities used in the approaches motivated from QCD. Applying these sum rules to the rho meson in nuclear medium with a simple pole-plus-continuum ansatz for the spectral densities, we found that the qualitative features of the QCD sum-rule predictions for the spectral parameters are consistent with these sum rules; however, the quantitative QCD sum-rule results violate the sum rules to certain degree.
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Study of hadronic properties at finite baryon density and temperature is essential in understanding the structure of matter. Recently, this subject has attracted much attention, motivated by the experimental attainment of hot and dense matter in heavy-ion collisions and by the theoretical expectations of a phase transition of matter from a hadronic phase to a quark-gluon plasma at high temperatures.

Despite the difficulties due to the nonperturbative features of QCD at large distances, one may study the properties of hadrons and the QCD vacuum by investigating the two-point correlation functions of currents, carrying the quantum numbers of the system under study. This approach is based on the analytic properties of the two-point correlation functions and on asymptotic freedom. The hadronic spectral properties (e.g., masses, coupling constants, etc.) appearing in the spectral densities can be related via dispersion relations to the correlation functions evaluated in terms of quark and gluon degrees of freedom. In practical applications such as QCD sum-rule method [1,2] and analyses of lattice QCD data [3] and interacting instanton model calculation [4], one needs to parametrize the spectral functions with a small number of parameters and to evaluate the correlation functions approximately [e.g., operator product expansion (OPE), lattice simulations, and interacting instanton approximation]. The success of such approaches depends on correct understanding of the qualitative features of the spectral functions and accurate evaluation of the correlation functions from QCD.

In this paper, we derive sum rules for the variation of finite-density spectral densities of vector channel with baryon density on the basis of dispersion relation and the OPE. These sum rules can be regarded as constraints on the phenomenological models for the finite-density spectral densities. We also apply these sum rules to the rho meson in nuclear medium with a simple pole-plus-continuum model for the spectral densities. We find that the qualitative features of the QCD sum-rule predictions for the spectral parameters are consistent with these constraints; the quantitative QCD sum-rule results, however, violate the constraints to certain degree.
Let us start with the correlation function of vector current at finite baryon density and zero temperature \([5,6]\):

\[
\Pi_{\mu\nu}(q, \rho_B) \equiv i \int d^4x e^{iq\cdot x} \langle T J_\mu(x) J_\nu(0) \rangle_{\rho_B} ,
\]

(1)

where \(J_\mu(x)\) is a vector current constructed from light quark fields [e.g., \(J_\mu(x) = \frac{1}{2} (\bar{u}(x) \gamma_\mu u(x) \pm \bar{d}(x) \gamma_\mu d(x))\)]. Throughout this paper, the up and down quark masses are taken to be equal and all the quark fields in the current are assumed to have the same mass. The notation \(\langle \cdots \rangle_{\rho_B}\) denotes the expectation value on the finite-density ground state characterized by the baryon density \(\rho_B\) in the rest frame and the four-velocity \(u^\mu\). In medium, there are in general two independent invariants in the vector channel, corresponding to the transverse and longitudinal polarizations \(\Pi_t(q, \rho_B)\) and \(\Pi_l(q, \rho_B)\). For simplicity, we will work in the rest frame of the medium, where \(u^\mu = (1, 0)\), and take the three momentum to be zero \(q = 0\). Then, since there is no specific spatial direction, the transverse polarization is related to the longitudinal one, \(\Pi_t(q_0, \rho_B) = q_0^2 \Pi_l(q_0, \rho_B)\), where \(\Pi_l(q_0, \rho_B) = \Pi_\mu(q_0, \rho_B)/(−3q_0^2)\) \([7,5,6]\). The longitudinal part, \(\Pi_l(q_0, \rho_B) = \Pi_l(q_0^2, \rho_B)\), satisfies the standard dispersion relation \([8,5]\)

\[
\Pi_l(q_0^2, \rho_B) = \int_0^\infty ds \rho(s, \rho_B) s + q_0^2 ,
\]

(2)

where \(\rho(s, \rho_B) = \pi^{-1} \text{Im} \Pi_l(s, \rho_B)\) is the finite-density spectral density. Here we have omitted the subtraction terms, which can be eliminated by taking derivatives of both sides of Eq. (2) with respect to \(q_0^2\). The sum rules to be derived are, however, independent of this process.

Using Eq. (2), one can write the difference of the correlation functions evaluated at different baryon densities as

\[
\Delta \Pi_l(Q^2) \equiv \Pi_l(Q^2, \rho_B) - \Pi_l(Q^2, \rho'_B) = \int_0^\infty ds \frac{\rho(s, \rho_B) - \rho(s, \rho'_B)}{s + Q^2} ,
\]

(3)

where \(Q^2 \equiv -q_0^2\) and \(\rho'_B\) denotes a different baryon density from \(\rho_B\).

At very short distances, or at very high energies, the difference between two correlation functions with different densities should go to zero due to asymptotic freedom of QCD. We
can now look for the consequences of this statement for the difference of spectral densities. Expanding the right-hand side of Eq. (3) for large values of $Q^2$ we get

\[ \Delta \Pi_l(Q^2) = \int_0^\infty ds \left[ \rho(s, \rho_B) - \rho(s, \rho'_B) \right] \left[ \frac{1}{Q^2} - \frac{s}{Q^4} + \frac{s^2}{Q^6} - \cdots \right]. \] (4)

On the other hand, for large $Q^2$ (i.e., in the deep Euclidean region), one can evaluate the correlation functions by expanding the product of currents according to the operator product expansion, which leads to

\[ \Pi_l(Q^2, \rho_B) = \sum_n C_n(Q^2) \langle \hat{O}_n \rangle_{\rho_B}, \] (5)

where $C_n(Q^2)$ are the Wilson coefficients and $\hat{O}_n$ are local composite operators constructed from quark and gluon fields. Here we have suppressed the dependence of both the coefficients and the operators on the normalization point $\mu$. The operators $\hat{O}_n$ are ordered by dimension (measured as a power of mass) and the $C_n(Q^2)$ for higher-dimensional operators fall off by corresponding powers of $Q^2$. The Wilson coefficients only depend on QCD Lagrangian parameters such as the quark masses and the strong coupling constant; all density dependence of the correlation function is included in the condensates $\langle \hat{O}_n \rangle_{\rho_B}$. Thus, one can express $\Delta \Pi_l(Q^2)$ as

\[ \Delta \Pi_l(Q^2) = \sum_n C_n(Q^2) \Delta \langle \hat{O}_n \rangle, \] (6)

where

\[ \Delta \langle \hat{O}_n \rangle \equiv \langle \hat{O}_n \rangle_{\rho_B} - \langle \hat{O}_n \rangle_{\rho'_B}. \] (7)

Note that the pure perturbative contribution [corresponding to the unit operator term, $\hat{O}_n = 1$, in the OPE] to the correlation function is independent of density, and thus does not appear in the difference $\Delta \Pi_l(Q^2)$. The lowest order contribution to $\Delta \Pi_l(Q^2)$ then comes from the condensates with lowest dimension ($\hat{O}_n \neq 1$), which, in the vector channel, are dimension four (including quark masses) condensates. Since $\Delta \Pi_l(Q^2)$ has dimension zero, the lowest order term in the OPE of $\Delta \Pi_l(Q^2)$ must be proportional to $1/Q^4$. However, the
lowest order term in the phenomenological representation Eq. (4) is proportional to $1/Q^2$. Therefore, we conclude that

$$\int_0^\infty [\rho(s, \rho_B) - \rho(s, \rho'_B)] \, ds = 0 . \quad (8)$$

This is a rigorous result. In the OPE framework it simply follows from the observation that the pure perturbative contribution is density blind and the lowest-dimensional condensates have dimension four. Although the finite-density spectral densities cannot be measured directly from experiments, Eq. (8) will constrain the change of the spectral density with baryon density. Phenomenological parametrizations often used in applications such as QCD sum-rule calculations or analyses of lattice QCD data, must satisfy this constraint. If one adopts a pole-plus-continuum ansatz for the spectral density, Eq. (8) indicates that the change in the coupling for the pole is equal to the shift in the continuum.

The OPE of $\Pi_l(Q^2, \rho_B)$ takes the general form

$$\Pi_l(Q^2) = \sum_i c_i^{(4)} \frac{\langle \hat{O}_i^{(4)} \rangle_{\rho_B}}{Q^4} + \sum_j c_j^{(6)} \frac{\langle \hat{O}_j^{(6)} \rangle_{\rho_B}}{Q^6} + \cdots , \quad (9)$$

where the ellipses denote the contributions of condensates with higher dimensions, the superscript indicates the dimension of operators, and the sum is over all contributing operators with a given dimension. The coefficients $c_i^{(d)}$ are dimensionless, and can, in principle, be dependent on $Q^2$ due to the perturbative corrections [only through strong coupling constant $\alpha_s(Q^2)$ and $m_q^2/Q^2$] [1]. For simplicity, we will work only to the lowest order in the strong coupling constant and to the first order in quark masses, where the coefficients $c_i^{(d)}$ become independent of $Q^2$.

We can then rewrite Eq. (3) as

$$\Delta \Pi_l(Q^2) = \sum_i c_i^{(4)} \frac{\Delta \langle \hat{O}_i^{(4)} \rangle}{Q^4} + \sum_j c_j^{(6)} \frac{\Delta \langle \hat{O}_j^{(6)} \rangle}{Q^6} + \cdots . \quad (10)$$

Comparing the coefficients of $1/Q^4$ in Eqs. (3) and (10), one obtains

$$\int_0^\infty s \left[ \rho(s, \rho_B) - \rho(s, \rho'_B) \right] \, ds = - \sum_i c_i^{(4)} \Delta \langle \hat{O}_i^{(4)} \rangle . \quad (11)$$
Similarly, equating the coefficients of $1/Q^6$, one finds

$$
\int_0^\infty s^2 [\rho(s, \rho_B) - \rho(s, \rho_B')] \, ds = \sum_i c_i^{(6)} \Delta(\hat{O}_i^{(6)}). \tag{12}
$$

Following the same pattern, one may derive an infinite series of sum rules, one for each OPE term (with fixed dimension) at small distances. The variation of the spectral density with baryon density must satisfy these sum rules. If one takes $\rho'_B = 0$ (i.e., in vacuum), these sum rules will constrain the change of the finite-density spectral density relative to the corresponding vacuum spectral density, which in some cases are experimentally accessible. In the QCD sum-rule applications or the analyses of lattice simulation data, one often parametrizes the spectral densities with a pole representing the lowest resonance plus continuum contribution roughly approximated by a perturbative evaluation of the correlation function, starting at an effective continuum threshold. One may apply the above sum rules to test this simple parametrization.

In principle, the sum rules described here can also be used to determine the finite-density spectral density provided that its corresponding vacuum spectral density and the values of in-medium and vacuum condensates are known. In practice, however, one has to truncate the OPE as the number of condensates with the same dimension appearing in higher order terms become larger and there is no reliable way to estimate these condensates. As a result, one can only expect to utilize the first few sum rules (at best), which, along with a simple ansatz for the spectral density, may give rise to an estimate of the finite-density spectral parameters.

We notice, however, that the sum rules of higher order (resulting from higher powers of $1/Q^2$) are more sensitive to the difference of the spectral density in higher energy region. In the QCD sum-rule approach, the Borel transformation suppresses the contribution from the higher energy region (continuum), though it introduces an auxiliary parameter (i.e., the Borel mass). Consequently, the spectral integral is saturated by the lowest resonance; the roughness of the approximation for the continuum is expected to have only small impact on
the spectral parameters for the lowest resonance \(^1\). The sum rules discussed here do not depend on any auxiliary parameter. However, for the sum rules to be useful in determining the spectral properties of hadrons, one needs to have a reliable model for the continuum.

The sum rules derived in the present paper may look like the usual finite energy sum rules \[^{3,8,10}\]. However, we emphasize that in our study, we focuses on the difference of the correlation function evaluated at two different baryon densities, instead of the correlation function at a particular density. The present sum rules are for the variation of the finite-density spectral densities with baryon density, instead of the spectral density at a particular density. Our derivation relies on the subtraction procedure and on the asymptotic freedom of QCD, which allows a short distance expansion of both the correlation function and the phenomenological dispersion relation. This technique has been used by Kapusta and Shuryak \[^{11}\] in deriving the Weinberg-type sum rules at zero and finite temperature. One may derive analogous sum rules by using other type of subtraction scheme, instead of different baryon density (see for example Ref. \[^{12}\]). One can also extend the sum rules to other channels, which will be documented in Ref. \[^{13}\].

We now turn to apply the sum rules Eqs. (8), (11) and (12) to the rho meson in nuclear medium, where the vector current is given by \(J_\mu(x) = \frac{1}{2} \left[ \bar{u}(x)\gamma_\mu u(x) - \bar{d}(x)\gamma_\mu d(x) \right] \). Since the complete spectral density in nuclear medium is not known experimentally, one cannot test directly whether the sum rules are indeed satisfied.

Various investigators \[^{3,8}\] have studied the properties of the rho meson in nuclear medium within QCD sum-rule approach. It is found that the rho meson mass (pole position of the propagator), the coupling of the vector current to the rho meson, and the effective continuum threshold all drop in nuclear medium. At nuclear matter saturation density, the rho meson mass decreases by \(\sim 15 - 18\%\) relative to its vacuum value. One can test whether these in-medium spectral features predicted from QCD sum-rule calculations are consistent with the sum rules derived in the present paper.

The explicit OPE result for the correlation function can be found in Refs. \[^{5,6}\]. To
the first order in the nucleon density $\rho_N$, the in-medium condensates can be written as 
\[ \langle \hat{O} \rangle_{\rho_N} = \langle \hat{O} \rangle_{\text{vac}} + \langle \hat{O} \rangle_{N} \rho_N, \]
where $\langle \hat{O} \rangle_{\text{vac}}$ is the spin averaged nucleon matrix element [5,14]. This linear approximation to the condensates is expected to be reasonable up to the nuclear matter saturation density [15,5].

Up to dimension six and to the linear order in $\rho_N$, the result for the difference $\Delta \Pi_l(Q^2)$ can be written as

\[ \Delta \Pi_l(Q^2) = \Pi_l(Q^2, \rho_N) - \Pi_l(Q^2, \rho_N') = 0 = \frac{m_q}{Q^4} \langle \bar{q} q \rangle_{N} \rho_N + \frac{1}{24Q^4} \left( \frac{\alpha_s}{\pi} G_{\mu\nu} G^{\mu\nu} \right)_N \rho_N \]

\[ + \frac{M_N}{4Q^4} A_{u+d}^2 \rho_N - \frac{224}{81Q^6} \pi \alpha_s \langle \bar{q} q \rangle_{\text{vac}} \langle \bar{q} q \rangle_{N} \rho_N - \frac{5M_N^3}{24Q^6} A_{u+d}^2 \rho_N \, , \]  

(13)

where $m_q$ is the average of up and down quark masses, $M_N$ is the nucleon mass and $A_{u+d}^2 (\equiv A^u_n + A^d_n)$ is a moment of the parton distributions in the deep inelastic scattering $A^q_n(\mu^2) = 2 \int_0^1 dx x^{n-1} [q(x, \mu^2) + (-1)^n \bar{q}(x, \mu^2)]$, where $q(x, \mu^2)$ and $\bar{q}(x, \mu^2)$ are the scale-dependent distribution functions for quarks and antiquarks (of flavor $q$) in the nucleon. Here the in-medium factorization approximation has been used for the four-quark condensates [i.e., $\langle \bar{q} q \rangle^2_{\rho_N} - \langle \bar{q} q \rangle^2_{\text{vac}} \simeq 2 \langle \bar{q} q \rangle_{\text{vac}} \langle \bar{q} q \rangle_{N} \rho_N$] [5]. In Eq. (13), we only retained the terms included in Refs. [3,4].

Since the sum rules Eqs. (8) and (11–12) involve integrations of the spectral densities with different powers of $s$, it is likely important to incorporate the finite widths of the resonances and to have a reliable model for the continuum. Here we will adopt the pole-plus-continuum parametrization for both vacuum and in-medium spectral density, which is widely used in the QCD sum-rule calculations [1,5,16], and hence only expect to test the qualitative features of the finite-density spectral parameters. The detailed account of the finite width of the rho meson, along with the incorporation of a more reliable continuum model, will be given in Ref. [13].

In the pole-plus-continuum approximation, one can write the vacuum spectral density as [1]

\[ \rho_{\text{vac}}(s) = F \delta(s - m^2_\rho) + \frac{1}{8\pi^2} \left( 1 + \frac{\alpha_s}{\pi} \right) \theta(s - s_0) \, , \]

(14)
and the finite-density spectral density as [3]

\[ \rho(s, \rho_N) = \rho_{sc} \delta(s) + F^* \delta(s - m^*_\rho) + \frac{1}{8\pi^2} \left( 1 + \frac{\alpha_s}{\pi} \right) \theta(s - s^*_0), \tag{15} \]

where \( F (F^*) \), \( m^*_\rho \) and \( s_0 (s^*_0) \) are the coupling, the rho meson mass and the continuum threshold, respectively, and \( \rho_{sc} \) denotes the contribution of Landau damping (or the scattering term) [7,5].

Substituting the spectral ansatz Eqs. (14–15) and the OPE results into Eqs. (8) and (11–12), we obtain the following sum rules

\[ \rho_{sc} + F^* - F + \frac{1}{8\pi^2} \left[ 1 + \frac{\alpha_s}{\pi} \right] (s_0 - s^*_0) = 0, \tag{16} \]

\[ F^* m^*_\rho^2 - F m^2_\rho + \frac{1}{16\pi^2} \left[ 1 + \frac{\alpha_s}{\pi} \right] (s^2_0 - s^2_0) = -m_q \langle \overline{q}q \rangle_N \rho_N - \frac{1}{24} (\frac{\alpha_s}{\pi} G_{\mu\nu}G^{\mu\nu})_N \rho_N - \frac{M_N}{4} A_{2}^{u+d} \rho_N, \tag{17} \]

\[ F^* m^4_\rho - F m^4_\rho + \frac{1}{24\pi^2} \left[ 1 + \frac{\alpha_s}{\pi} \right] (s^3_0 - s^3_0) = -\frac{224}{81} \pi \alpha_s \langle \overline{q}q \rangle_{\text{vac}} \langle \overline{q}q \rangle_N \rho_N - \frac{5M_N^3}{24} A_{4}^{u+d} \rho_N, \tag{18} \]

where we have assumed that the in-medium continuum threshold \( s^*_0 \) is less than its vacuum value \( s_0 \). Substituting the QCD sum-rule predictions for the spectral parameters into these sum rules one may check how well these sum rules are satisfied. Alternatively, one may extract the in-medium spectral parameters by solving these three equations and compare the results with the QCD sum-rule predictions. Here we follow the latter.

To obtain the finite-density spectral parameters \( m^*_\rho \), \( F^* \), and \( s^*_0 \) from Eqs. (16–18), one needs to know the various nucleon matrix elements appearing on the right hand sides as well as the corresponding vacuum spectral parameters. The nucleon matrix element \( \langle \overline{q}q \rangle_N \) is related to the nucleon sigma term \( \langle \overline{q}q \rangle_N = \sigma_N / 2m_q \); we take \( \sigma_N \approx 45 \text{ MeV} \) [17] and \( m_q \approx 5.5 \text{ MeV} \) [13,14]. For the gluon matrix element, we use \( \langle (\alpha_s / \pi) G_{\mu\nu}G^{\mu\nu} \rangle_N \approx -650 \text{ MeV} \) [15,14]. The moments of parton distribution are taken to be \( A_{2}^{u+d} \approx 0.938 \) and \( A_{4}^{u+d} \approx 0.121 \) (at \( \mu^2 = 1 \text{ GeV}^2 \)) [3]. We adopt \( \langle \overline{q}q \rangle_{\text{vac}} \approx (-245 \text{ MeV})^3 \) [18] and \( \alpha_s \approx 0.3 \) [3] in our
calculations. The nuclear matter saturation density is taken to be \( \rho_0 = (110 \text{ MeV})^3 \). We fix \( m_\rho = 770 \text{ MeV} \) and \( s_0 = 1.5 \text{ GeV}^2 \) [1], and parametrize the scattering term as \( \rho_{sc} = a_0 \rho_N \).

In Fig. 1, the resulting ratio of the in-medium rho meson mass to its vacuum value is plotted as function of nucleon density for different values of \( a_0 \) and a fixed \( F \) value \( F = 2 f_\pi^2 \) with \( f_\pi = 93.5 \text{ MeV} \), which is obtained by using \( F = m_\rho^2/g_\rho^2 \) [13] with the KSFR relation \( g_\rho^2 = m_\rho^2/2 f_\pi^2 \). One notices that the in-medium rho meson mass drops relative to its vacuum value. Similar behavior is also found for both the coupling and the continuum threshold. These qualitative features are consistent with those predicted from the QCD sum-rule calculations.

On the other hand, we observe that the quantitative result for the ratio is sensitive to the value of \( a_0 \) (i.e., \( \rho_{sc} \)). For \( a_0 = 0 \), the in-medium rho meson mass is only few percent smaller than its free space value even at the nuclear matter density. For \( a_0 = 3.6 \text{ GeV}^{-1} \), the rho meson mass drops to \( \sim 0.8 m_\rho \) at the saturation density. In Ref. [3] the scattering term is neglected (i.e., \( a_0 = 0 \)) while in Ref. [4] it is taken to be \( a_0 = 1/2 M_N \). The ratio \( m_\rho^*/m_\rho \) is found to be \( \sim 15 - 18\% \) at \( \rho_N = \rho_0 \) in these references. However, we find that to obtain a \( 15 - 18\% \) decrease in the rho mass at the nuclear matter saturation density, one needs to use a value of \( a_0 \sim 3 \text{ GeV}^{-1} \), which is much larger than that used in Ref. [4]. This inconsistency signals that the simple pole-plus-continuum model for the spectral densities with the QCD sum-rule predictions for the spectral parameters violates the constraints Eqs. (8) and (11–12) to certain degree. A improved model beyond this simple parametrization for the spectral densities may be needed to satisfy the constraints.

In Fig. 2 we plot the ratio \( m_\rho^*/m_\rho \) as a function of the nucleon density for different values of \( F \) with fixed \( a_0 = 1/2 M_N \). It can be seen that the ratio is very sensitive to \( F \), in particular for small \( F \) values. (For \( F \geq 0.035 \text{ GeV}^2 \), there is no real solution). Again, we note that a much smaller value of \( F \) is necessary to reproduce the QCD sum-rule result. It is also found that our result is relatively insensitive to the values of \( m_\rho \) and \( s_0 \).

In conclusion, we have derived sum rules for the variation of the finite-density spectral
density of vector channel with baryon density within the framework of operator product expansion and the dispersion relation. These sum rules may serve as constraints on the phenomenological models used in the QCD sum-rule calculations or in the interpretation of the lattice QCD data. We also noted that in principle one can use these sum rules to determine the qualitative properties of the finite-density spectral parameters if the corresponding vacuum spectral density and the in-medium and vacuum condensates are known. We applied the first three sum rules to the rho meson in nuclear medium with a simple pole-plus-continuum parametrization for the spectral densities, and found that the qualitative features of the QCD sum-rule predictions are consistent with our sum rules. However, the quantitative result shows that the simple ansatz with QCD sum-rule predictions for the spectral parameters violates the sum rules to certain degree.

This suggests that the inclusion of the finite widths of the resonances and the refinement of the continuum model may be important. This point, along with the full detail of the present paper and its extension to other channels, will be reported elsewhere [13].

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FIGURES

FIG. 1. Ratio $m^*/m$, obtained from solving Eqs. (16–18), as a function of the nucleon density. The four curves correspond to $a_0 = 0$ (solid), $a_0 = 1.2$ GeV$^{-1}$ (long-dashed), $a_0 = 2.4$ GeV$^{-1}$ (dotted), and $a_0 = 3.6$ GeV$^{-1}$ (dashed). The other input parameters are described in the text.

FIG. 2. Ratio $m^*/m$, obtained from solving Eqs. (16–18), as a function of the nucleon density. The four curves correspond to $F = 0.005$ GeV$^2$ (solid), $F = 0.01$ GeV$^2$ (long-dashed), $F = 0.02$ GeV$^2$ (dotted), and $F = 0.03$ GeV$^2$ (dashed). The other input parameters are the same as in Fig. 1.
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