Electromagnetic form factors of the ρ meson in light cone QCD sum rules

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Abstract

We investigate the electromagnetic form factors of the ρ meson in light cone QCD sum rules. We find that the ratio of the magnetic and charge form factors is larger than two at all values of Q^2, ($Q^2 \geq 0.5\ GeV^2$). The values of the individual form factors at fixed values of Q^2 predicted by the light cone QCD sum rules are quite different compared to the results of other approaches. These results can be checked in future, when more precise data on ρ meson form factors is available.

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1 Introduction

The QCD sum rules method [1] is one of the most powerful tools in studying low energy hadron physics. In this method, physically measurable quantities of hadrons are connected with QCD parameters, where hadrons are represented by their interpolating quark currents. The main idea of this method is to calculate the correlator functions of the interpolating quark currents in the deep Eucledian region with the help of operator product expansion (OPE) which allows one to take into account both perturbative and nonperturbative contributions. Relevant physical quantities are determined by matching the correlator to its phenomenological representation.

In the current literature, a new, widely discussed alternative to traditional sum rules, namely, QCD light cone sum rules (QLCSR) is a convenient tool for the study of exclusive processes. This method is based on OPE on the light cone, which is an expansion over twists of the operators, rather than dimensions, as is the case in the traditional QCD sum rules. Moreover, in this method, all the nonperturbative dynamics encoded in the light cone distribution amplitudes, determine the matrix elements of the nonlocal operators between the vacuum and and the hadronic states (more about this method and its applications can be found in [2, 3]).

In the present work, we study the ρ meson form factors in QLCSR. It should be mentioned here that the ρ meson form factors are calculated at intermediate momentum transfer by using the three–point QCD sum rules method in [4]. Recently, ρ meson form factors, including next–to–leading–order perturbation theory, are analyzed within the same framework [5]. The $Q^2 = 0$ point lies outside the applicability region of the three–point QCD sum rules. Consequently, extrapolating these form factors to $Q^2 = 0$, in principle, one can determine static characteristics of the ρ meson, such as magnetic and quadrapole moments. The ρ meson magnetic moment has already been investigated in the framework of the traditional three–point and QLCSR in [6] and [7], respectively. It should be noted here that, the QLCSR is successfully applied to a wide range of problems of the hadron physics, for example, magnetic moments of the octet and decuplet baryons are calculated in [8] and [9], and magnetic moment of the nucleon in [10], respectively.

The paper is organized as follows. In section 2, we derive the sum rules for the ρ meson electromagnetic form factors. In section 3, our numerical results and a comparison of them with the results of various approaches existing in the literature is presented.

2 Theoretical framework

In order to determine ρ meson electromagnetic form factors we will use QLCSR method. For this purpose, we consider the following correlator function

$$\Pi_{\mu\nu}(p,q) = i \int d^4x e^{iqx} \langle \rho(p) | J_\mu^{el}(x) J_\nu(0) | 0 \rangle \ , \qquad (1)$$

where

$$\begin{aligned} J_\nu(x) &= \bar{u}(x)\gamma_\nu d(x) \ , \\ J_\mu^{el}(x) &= \sum_{q=u,d} e_q \bar{q}(x)\gamma_\mu q(x) \ , \end{aligned}$$

are the ρ meson interpolating and electromagnetic currents, respectively.

Firstly, we calculate the phenomenological part of the correlator. By inserting the complete set of states which have the quantum numbers of the ρ meson, between the currents in Eq. (1), we get

$$\Pi_{\mu\nu} = \frac{\langle \rho(p,\varepsilon)|J_\mu^{el}(x)|\rho(p',\varepsilon')\rangle \langle \rho(p',\varepsilon')|J_\nu(0)|0\rangle}{p'^2 - m_\rho^2} + \cdots \ , \tag{2}$$

where $p' = p + q$, q is the momentum of the electromagnetic current, ε is the polarization of the ρ meson and $\cdots$ describes contributions from higher states. The matrix element $\langle \rho|j_\nu|0\rangle$ is determined as

$$\langle \rho|j_\nu|0\rangle = f_\rho m_\rho \varepsilon_\nu^*(p) \ . \tag{3}$$

Assuming parity and time–reversal invariance, the electromagnetic vertex of the ρ meson can be written in terms of three Lorentz invariant form factors [11]

$$\begin{aligned}\langle \rho(p,\varepsilon)|J_\mu^{el}|\rho(\rho',\varepsilon')\rangle = \varepsilon^{*\alpha}\varepsilon'^{\beta}\Big\{ &- G_1(Q^2) g_{\alpha\beta}(p+p')_\mu - G_2(Q^2)(q_\beta g_{\mu\alpha} - q_\alpha g_{\mu\beta}) \\ &+ \frac{1}{2m_\rho^2} G_3(Q^2) q_\alpha q_\beta (p+p')_\mu \Big\} \ ,\end{aligned} \tag{4}$$

where $Q^2 = -q^2$ is the square of the momentum transfer. It should be mentioned here that, in practical computation, instead of calculating the Lorentz invariant form factors $G_i(Q^2)$, the physical charge G_C, the magnetic G_M and quadrapole G_Q form factors are often used.

The Lorentz invariant form factors $G_i(Q^2)$ are related to the charge, magnetic and quadrapole form factors through the following relations

$$\begin{aligned} G_C &= G_1 + \frac{2}{3}\eta F_{\mathcal{D}} \ , \\ G_M &= G_2 \ , \\ G_Q &= G_1 - G_2 + (1+\eta) G_3 \ , \end{aligned} \tag{5}$$

where $\eta = Q^2/4m_\rho^2$. At zero momentum transfer, these form factors are proportional to the usual static quantities of charge e, magnetic moment μ and quadrapole moment $\mathcal{D}$

$$\begin{aligned} eG_C(0) &= e \ , \\ eG_M(0) &= 2m_\rho \mu \ , \\ eG_Q(0) &= m_\rho^2 \mathcal{D} \ . \end{aligned} \tag{6}$$

Substituting Eqs. (3) and (4) into Eq, (2), the phenomenological part of the correlator takes the following form

$$\begin{aligned}\Pi_{\mu\nu} = \frac{f_\rho m_\rho}{p'^2 - m_\rho^2}\Bigg\{ & G_1 (p+p')_\mu \left[\varepsilon_\nu - \frac{(\varepsilon q) p'_\nu}{m_\rho^2}\right] + G_2 \Bigg[\varepsilon_\mu q_\nu - (\varepsilon q) g_{\mu\nu} \\ &+ \frac{1}{m_\rho^2}\left(\frac{1}{2} Q^2 \varepsilon_\mu p'_\nu + (q\varepsilon) p'_\mu p'_\nu\right)\Bigg] - G_3 \frac{(p+p')_\mu}{2m_\rho^2}(q\varepsilon)\left(q_\nu + \frac{Q^2}{2m_\rho^2} p'_\nu\right)\Bigg\} \ .\end{aligned} \tag{7}$$

It is apparent from Eq. (7) that there are many different structures each of which can be used to extract the above–mentioned form factors.

In order to find out the ρ meson electromagnetic form factors, we pick the following three structures: $\varepsilon_\nu p_\mu$, $\varepsilon_\mu q_\nu$ and $(\varepsilon q) p_\mu q_\nu$. Hence we can write

$$\Pi_{\mu\nu}(p,q) = \Pi_1(p,q)\varepsilon_\nu p_\mu + \Pi_2(p,q)\varepsilon_\mu q_\nu + \Pi_3(p,q)(\varepsilon q)p_\mu q_\nu + \cdots ~, \tag{8}$$

where

$$\begin{aligned}
\Pi_1(p,q) &= -G_1(Q^2)\frac{2 f_\rho m_\rho}{m_\rho^2 - (p+q)^2} ~, \\
\Pi_2(p,q) &= -G_2(Q^2)(1+2\eta)\frac{f_\rho m_\rho}{m_\rho^2 - (p+q)^2} ~, \\
\Pi_3(p,q) &= [2G_1(Q^2) - G_2(Q^2) + G_3(Q^2)(1+2\eta)]\frac{f_\rho m_\rho}{m_\rho^2[m_\rho^2 - (p+q)^2]} ~.
\end{aligned} \tag{9}$$

For the invariant amplitudes $\Pi_i(p,q)$, one can write a general dispersion relation in $(p+q)^2$ in the form

$$\Pi_i(p,q) = \int ds \frac{\rho_i(s,q)}{s-(p+q)^2} + \text{subtr.} ~, \tag{10}$$

where the spectral density corresponding to Eq. (7) is

$$\rho_i(s,q) = A F_i(Q^2)\delta(s-m_\rho^2) ~, \tag{11}$$

where

$$\begin{aligned}
F_1 &= -2G_1(Q^2) f_\rho m_\rho ~, \\
F_2 &= -G_2(Q^2)(1+2\eta) f_\rho m_\rho ~, \\
F_3 &= [2G_1(Q^2) - G_2(Q^2) + G_3(Q^2)(1+2\eta)]\frac{f_\rho m_\rho}{m_\rho^2} ~.
\end{aligned}$$

According to QCD sum rules philosophy, in constructing sum rules for the form factors we need representation of the correlator function from QCD side. After contracting quark fields, the correlator function takes the form:

$$\Pi_{\mu\nu} = \int d^4x e^{iqx} \langle \rho(p) | e_u \bar{u}(x)\gamma_\mu S_d(x,0)\gamma_\nu d(0) + e_d \bar{u}(0)\gamma_\nu S_u(0,x)\gamma_\mu d(x) | 0 \rangle ~, \tag{12}$$

where $S_q(x)$ is the full quark operator of u (or d) quark. Imposing $SU(2)$ symmetry and neglecting u and d quark masses, we have $S_u(x,0) = S_d(x,0) = S(x)$, and hence the full light quark propagator takes the form

$$S(x) = \frac{i \not{x}}{2\pi^2 x^4} - \frac{\langle \bar{q} q \rangle}{12} - \frac{x^2 m_0^2}{192}\langle \bar{q} q \rangle - i g_s \int_0^1 dv \left[\frac{\not{x}\sigma_{\alpha\beta}}{16\pi^2 x^2} G_{\alpha\beta}(vx) - \frac{i}{4\pi^2 x^2} v x_\alpha G_{\alpha\beta}\gamma_\beta \right] \tag{13}$$

Note that the second and third terms of $S(x)$ do not give any contribution to the considered problem after Borel transformation is carried out (see below). Rewriting Eq. (13) in momentum representation and substituting it in Eq. (12), we get

$$\begin{aligned}
\Pi_{\mu\nu} = \frac{i^2}{4}\int d^4x d^4k e^{i(q-k)x}\Big\{ & e_u \langle \rho(p)|\bar{u}(x)\Gamma_i d(0)|0\rangle \mathrm{Tr}\gamma_\mu \frac{\not k}{k^2}\gamma_\nu\Gamma_i \\
- & e_u \langle \rho(p)|\bar{u}(x) g G_{\alpha\beta}(vx)\Gamma_i d(0)|0\rangle \mathrm{Tr}\gamma_\mu \left[\frac{1}{2}\frac{\not k}{k^4}\sigma_{\alpha\beta} - \frac{v}{k^2}x_\alpha\gamma_\beta\right]\gamma_\nu\Gamma_i \\
+ & e_d \langle \rho(p)|\bar{u}(0)\Gamma_i d(x)|0\rangle \mathrm{Tr}\gamma_\nu \frac{\not k}{k^2}\gamma_\mu\Gamma_i \\
- & e_d \langle \rho(p)|\bar{u}(0) g G_{\alpha\beta}(vx)\Gamma_i d(x)|0\rangle \mathrm{Tr}\gamma_\nu \left[\frac{1}{2}\frac{\not k}{k^4}\sigma_{\alpha\beta} - \frac{v}{k^2}x_\alpha\gamma_\beta\right]\gamma_\mu\Gamma_i \Big\} ~, \qquad (14)
\end{aligned}$$

where $\Gamma_i = (I, \gamma_5, \gamma_\alpha, \gamma_\alpha\gamma_5, \sigma_{\alpha\beta})$ is the full set of the Dirac matrices. It follows from this expression that, in order to calculate the correlator from QCD side one needs to know the matrix element of the nonlocal operators between vacuum and the ρ meson, i.e., ρ meson distribution amplitudes (or wave functions). We can see from Eq. (14) that main contribution to the correlator comes from only from the wave functions that contain odd–number of γ–matrices that are defined in the following way [14]:

$$\begin{aligned}
\langle \rho(p,\varepsilon)|\bar{u}(x)\gamma_\mu d(0)|0\rangle = f_\rho m_\rho \Big\{ & \frac{\varepsilon x}{px}\int_0^1 dv e^{ivpx}\left[\phi_\parallel(v,\mu) + \frac{m_\rho^2 x^2}{16}A(v,\mu)\right] p_\mu \\
+ & \left(\varepsilon_\mu - \frac{\varepsilon x}{px}p_\mu\right)\int_0^1 dv e^{ivpx} g_\perp^v(v,x) \\
- & \frac{1}{2}x_\mu \frac{\varepsilon x}{(px)^2} m_\rho^2 \int_0^1 dv e^{ivpx} C(v,\mu) \Big\} ~, \qquad (15)
\end{aligned}$$

$$\langle \rho(p,\varepsilon)|\bar{u}(x)\gamma_\mu\gamma_5 d(0)|0\rangle = \frac{1}{4} f_\rho m_\rho \epsilon_{\mu\alpha\beta\delta}\varepsilon_\alpha p_\beta x_\delta \int_0^1 dv e^{ivpx} g_\perp^a(v,\mu) ~, \qquad (16)$$

$$\begin{aligned}
\langle \rho(p,\varepsilon)|\bar{u}(x) g G_{\mu\nu}\gamma_\alpha d(0)|0\rangle = & -i f_\rho m_\rho P_\alpha \Big(P_\nu \varepsilon_\mu^\perp - P_\mu \varepsilon_\nu^\perp\Big)\mathcal{V} \\
& - i f_\rho m_\rho^3 \frac{\varepsilon x}{px}\Big(P_\mu g_{\alpha\nu}^\perp - P_\nu g_{\alpha\mu}^\perp\Big)\Phi \\
& - i f_\rho m_\rho^3 \frac{\varepsilon x}{(px)^2} P_\alpha \Big(P_\mu x_\nu - P_\nu x_\mu\Big)\Psi ~, \qquad (17)
\end{aligned}$$

$$\begin{aligned}
\langle \rho(p,\varepsilon)|\bar{u}(x) g \widetilde{G}_{\mu\nu}\gamma_\alpha\gamma_5 d(0)|0\rangle = & f_\rho m_\rho P_\alpha \Big(P_\nu \varepsilon_\mu^\perp - P_\mu \varepsilon_\nu^\perp\Big) A \\
& + i f_\rho m_\rho^3 \frac{\varepsilon x}{px}\Big(P_\mu g_{\alpha\nu}^\perp - P_\nu g_{\alpha\mu}^\perp\Big)\widetilde{\Phi} \\
& + f_\rho m_\rho^3 \frac{\varepsilon x}{(px)^2} P_\alpha \Big(P_\mu x_\nu - P_\nu x_\mu\Big)\widetilde{\Psi} ~, \qquad (18)
\end{aligned}$$

where P_μ, $\varepsilon_\mu^\perp$ and $g_{\mu\nu}^\perp$ are defined as

$$P_\mu = p_\mu - \frac{m_\rho^2}{2px}x_\mu ~,$$

$$\varepsilon_\mu^\perp = \varepsilon_\mu - \frac{\varepsilon x}{px}\left(P_\mu - \frac{m_\rho^2}{px} x_\mu\right) ,$$
$$g_{\mu\nu}^\perp = \left(g_{\mu\nu} - \frac{P_\mu x_\nu + P_\nu x_\nu}{px}\right) ,$$

where $px = Px$ has been used.

In Eqs. (15)–(18), $\phi_\parallel(v,\mu)$ is the leading twist–2 wave function, while $g_\perp^v$, $g_\perp^a$ and $\mathcal{V}$ are the twist–3 and all the remaining ones are the twist–4 wave functions. In Eqs. (17) and (18) the following relation is used

$$\mathcal{V}(v,px) = \int \mathcal{D}\alpha e^{ipx(\alpha_1+v\alpha_3)} \mathcal{V}(\alpha_1,\alpha_2,\alpha_3) ,$$

where

$$\mathcal{D}\alpha = d\alpha_1 d\alpha_2 d\alpha_3 \delta(1-\alpha_1-\alpha_2-\alpha_3) .$$

In order to suppress the contributions of higher states and continuum and further eliminate the subtraction terms in the dispersion relation, it is necessary to perform Borel transformation with respect to $-(p+q)^2$. The contributions of higher states and continuum are subtracted using the quark hadron duality (for details, see [2, 3], [8]–[10], [12] and [13]).

The sum rules for form factors can be obtained after applying Borel transformation to the two different representation of the invariant functions F_i and then matching these results, and in doing so, we get

$$\begin{aligned}
G_1(Q^2) &= \frac{1}{2} e^{m_\rho^2/M^2} \Bigg\{ \int_{v_0}^1 dv e^{-s(v)/M^2} \Bigg[-\frac{m_\rho^2}{4} \frac{1}{v^2 M^2} \Big(e_u A^i(v) - e_d A^i(\bar v) \Big) \\
&+ \frac{1}{v} \Big(e_u \Phi^i(v) - e_d \Phi^i(\bar v) \Big) - \frac{2m_\rho^2}{vM^2} \Big(e_u C^{ii}(v) - e_d C^{ii}(\bar v) \Big) \\
&+ \Big(e_u g_\perp^v(v) - e_d g_\perp^v(\bar v) \Big) + \frac{1}{4} \left(\frac{1}{v} - \frac{Q^2 + m_\rho^2 v^2}{v^2 M^2} \right) \Big(e_u g_\perp^a(v) + e_d g_\perp^a(\bar v) \Big) \Bigg] \Bigg\} , \qquad (19)
\end{aligned}$$

$$\begin{aligned}
G_2(Q^2) &= \frac{2m_\rho^2}{Q^2 + 2m_\rho^2} e^{m_\rho^2/M^2} \Bigg\{ \int_{v_0}^1 dv e^{-s(v)/M^2} \Bigg[-\frac{2m_\rho^2}{v^2 M^2} \Big(e_u C^{ii}(v) - e_d C^{ii}(\bar v) \Big) \\
&+ \frac{1}{v} \Big(e_u g_\perp^v(v) - e_d g_\perp^v(\bar v) \Big) - \frac{1}{4v^2} \left(1 - \frac{Q^2 + m_\rho^2 v^2}{vM^2} \right) \Big(e_u g_\perp^a(v) + e_d g_\perp^a(\bar v) \Big) \Bigg] \Bigg\} , \qquad (20)
\end{aligned}$$

$$\begin{aligned}
-2G_1 + G_2 - G_3(1+2\eta) &= m_\rho^2 e^{m_\rho^2/M^2} \Bigg\{ \int_{v_0}^1 dv e^{-s(v)/M^2} \Bigg[-\frac{m_\rho^2}{2v^3 M^4} \Big(e_u A^i(v) \\
&- e_d A^i(\bar v) \Big) - \frac{1}{2v^2 M^2} \Big(e_u g_\perp^a(v) + e_d g_\perp^a(\bar v) \Big) + \frac{2}{v^2 M^2} \Big(e_u \Phi^i(v) - e_d \Phi^i(\bar v) \Big) \\
&- \frac{4m_\rho^2 v}{v^3 M^4} \Big(e_u C^{ii}(v) - e_d C^{ii}(\bar v) \Big) \Bigg] \Bigg\} , \qquad (21)
\end{aligned}$$

where M^2 is the Borel parameter and

$$s(v) = \frac{(Q^2 + m_\rho^2 v)\bar v}{v} ,$$

$$\bar{v} = 1 - v \ ,$$
$$v_0 = -\frac{1}{2m_\rho^2}\Big[(s_0 + Q^2 - m_\rho^2)^2 - \sqrt{(s_0 + Q^2 - m_\rho^2)^2 + 4m_\rho^2 Q^2}\Big] \ ,$$

where s_0 is the continuum threshold.

In Eqs. (19)–(21), functions A^i, Φ^i and C^{ii} are defined as follows

$$A^i = -\int_0^v duA(u) \ ,$$
$$\Phi^i(v) = -\int_0^v du\Phi(u) \ ,$$
$$C^i(v) = -\int_0^v duC(u) \ ,$$
$$C^{ii}(v) = -\int_0^v duC^i(u) \ .$$

Note that the terms involving three–particle ρ meson wave functions are not presented in Eqs. (19)–(21) since their expressions are rather lengthy, however, their contributions are taken into account in the numerical analysis, which constitute about 5% of the total result.

3 Numerical analysis

In this section we present our numerical calculations on charge, magnetic and quadrapole form factors. The main input parameters of the QLCSR in regard to the above–mentioned form factors are the ρ meson wave functions, whose explicit expressions are given in [12] and [14], and we use them in our analysis.

Apart from the wave functions, the sum rules for the form factors depend on the value of the continuum threshold s_0 and Borel parameter M^2. In the present work, we calculate the form factors at three different values of the threshold, i.e., $s_0 = 1.8\ GeV^2$, $s_0 = 2.0\ GeV^2$ and $s_0 = 2.2\ GeV^2$ (see also [4]).

In Figs. (1), (2) and (3) we present the dependence of G_C, G_M and $G_{\mathcal{D}}$ on M^2 at six different values $Q^2 = 0.5\ GeV^2$, $Q^2 = 1.0\ GeV^2$, $Q^2 = 2.0\ GeV^2$, $Q^2 = 3.0\ GeV^2$, $Q^2 = 4.0\ GeV^2$ and $Q^2 = 5.0\ GeV^2$ of the momentum transfer, respectively. The Borel parameter M^2 in the sum rule is an auxiliary parameter and therefore physical quantities must be independent of it. For this reason we must determine the region of M^2 in which the form factors are independent of its value. In determining the working region of of the Borel parameter M^2, the following two conditions must be satisfied:

- M^2 should be large enough to suppress the contributions coming from higher twists, and,
- M^2 should be small enough in order to suppress the continuum and higher state contributions.

In the present analysis, both conditions are satisfied for the three form factors when M^2 varies in the region $1.0\ GeV^2 \le M^2 \le 2.5\ GeV^2$.

In Figs. (4), (5) and (6) we present the dependence of the charge G_C, magnetic G_M and quadrapole G_Q form factors on Q^2, at three different values $s_0 = 1.8~GeV^2$, $s_0 = 2.0~GeV^2$ and $s_0 = 2.2~GeV^2$ of the threshold, and at the fixed value $M^2 = 1.0~GeV^2$ of the Borel parameter.

We observe from these figures that the dependence of the form factors on s_0 is rather weak and when s_0 vary from $s_0 = 1.8~GeV^2$ to $s_0 = 2.2~GeV^2$. Unfortunately, the sum rules fail working at small Q^2, and hence do not allow determination of the magnetic moment of the ρ meson with better accuracy compared to the results in the literature, such as, the results predicted by the sum rules [6, 7], Dyson–Schwinger based models [15, 16], Covariant light–front approach with constituent quark model [17], light–front formalism [18] and light–front quark model [19]. Our analysis predicts that, starting from $Q^2 = 0.5~GeV^2$, the ratio $G_M(Q^2)/G_C(Q^2)$ is around 2.3 at all values of Q^2. If we assume that this behavior holds at smaller values of Q^2, we can conclude that $\mu \simeq 2.3$, which is quite close to the prediction on the magnetic moment of the ρ meson [7]. It should be noted here that the values of the form factors at different values of Q^2 in this work are quite different compared to the predictions of the other approaches. For example, our results for $G_C(Q^2)$ and $G_M(Q^2)$ at all values of Q^2 are, approximately, two times smaller, while $G_Q(Q^2)$ is two times larger (in magnitude) compared to the prediction of [19] (we use the same parametrization for the form factors as in [19]). Therefore, more data on ρ meson is needed in order to choose the "right" model in calculating the form factors.

In conclusion, we have presented the results for the ρ meson form factors in the frame work of QLCSR. We have obtained that, in the region of applicability of the method, at all values of Q^2, the ratio $G_M(Q^2)/G_C(Q^2)$ is larger than 2, more precisely, this ratio varies around 2.3. This result can be checked when more precise data on ρ meson form factors is available.

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Figure captions

Fig. (1) The dependence of the charge form factor G_C on M^2, at six fixed values of Q^2, and at $s_0 = 2.2\ GeV^2$.

Fig. (2) The same as in Fig. (1), but for the magnetic form factor G_M.

Fig. (3) The same as in Fig. (1), but for the quadrapole form factor G_Q.

Fig. (4) The dependence of the charge form factor G_C on Q^2, at Three fixed values of s_0, and at $M^2 = 1.0\ GeV^2$.

Fig. (5) The same as in Fig. (4), but for the magnetic form factor G_M.

Fig. (6) The same as in Fig. (4), but for the quadrapole form factor G_Q.

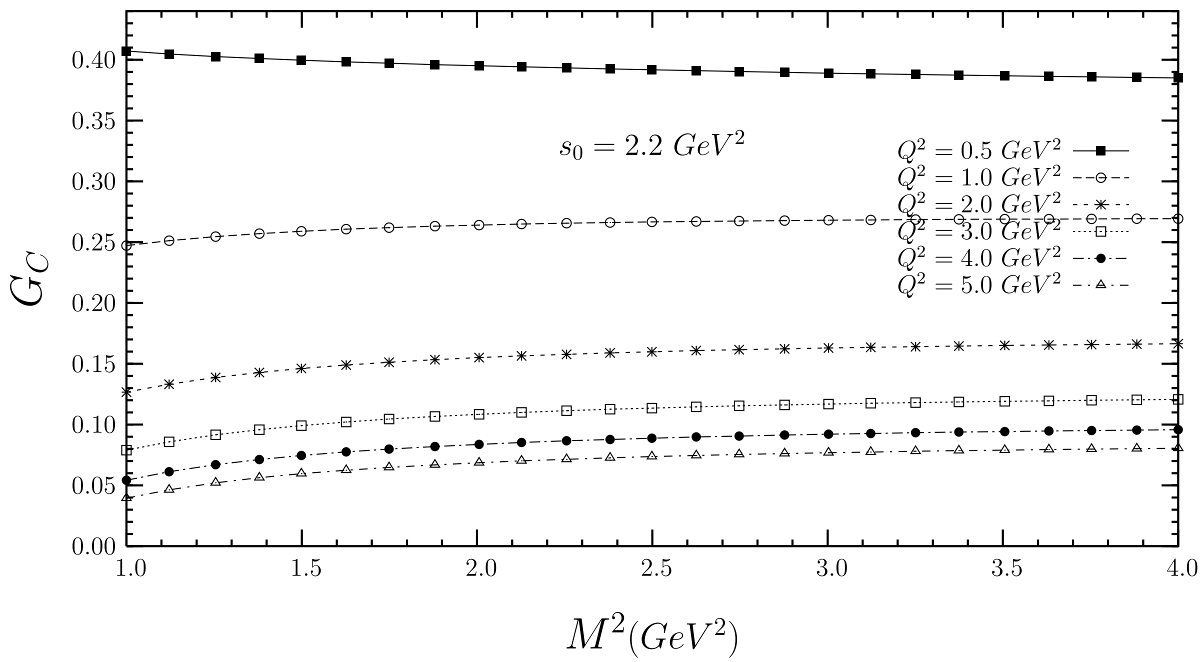
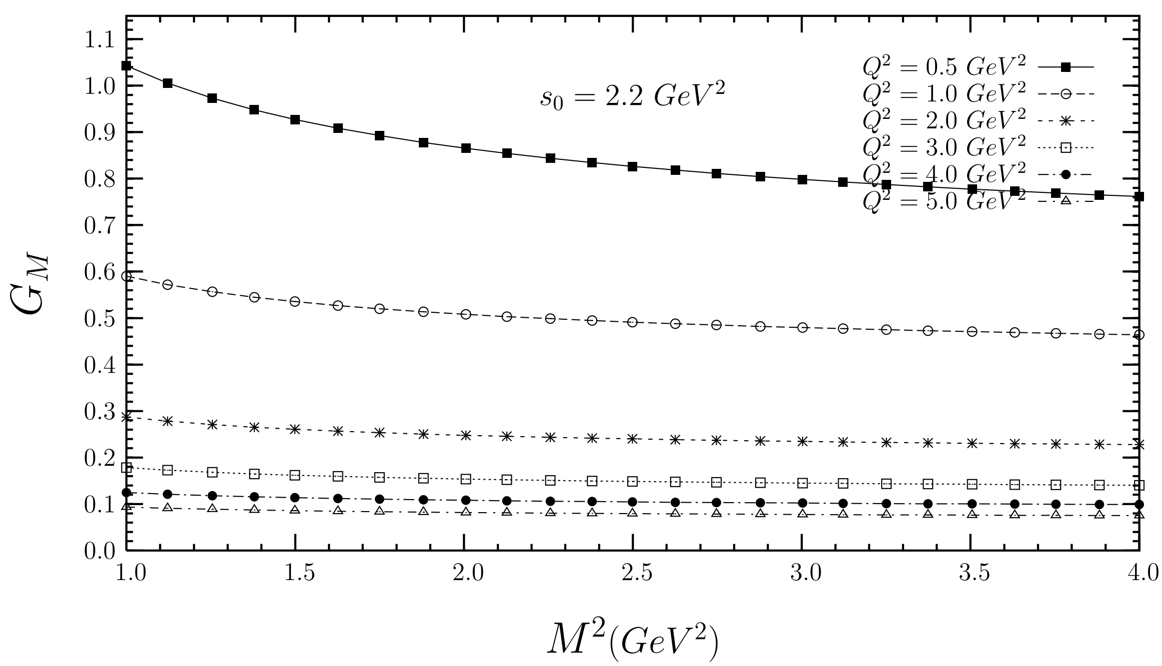


Figure 1:



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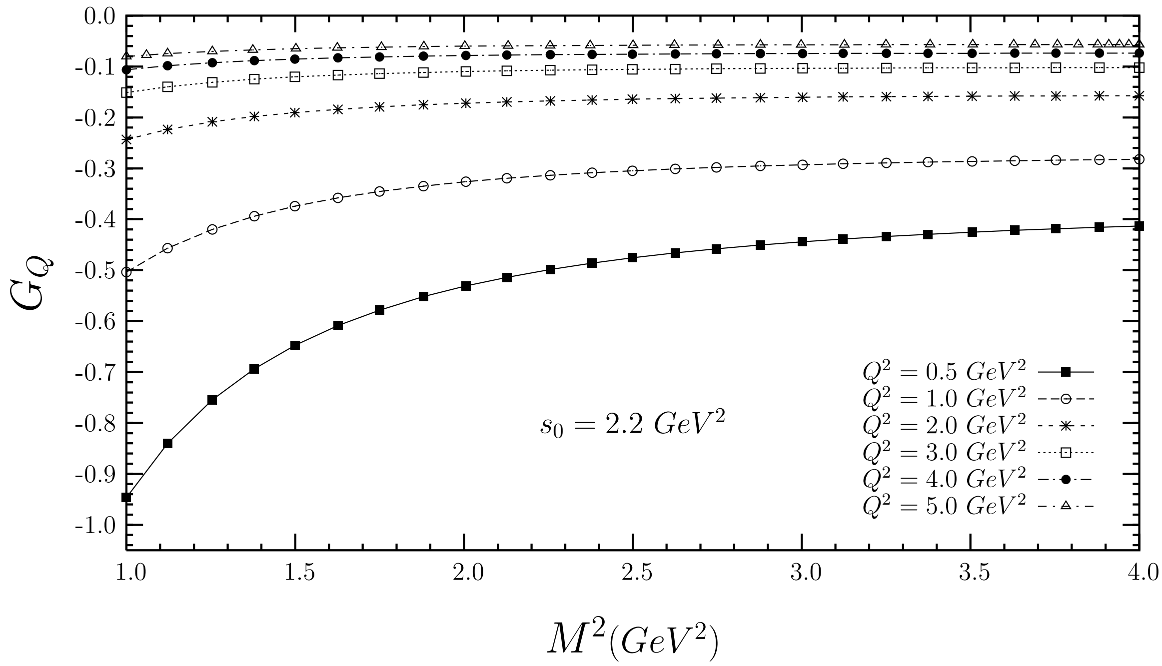

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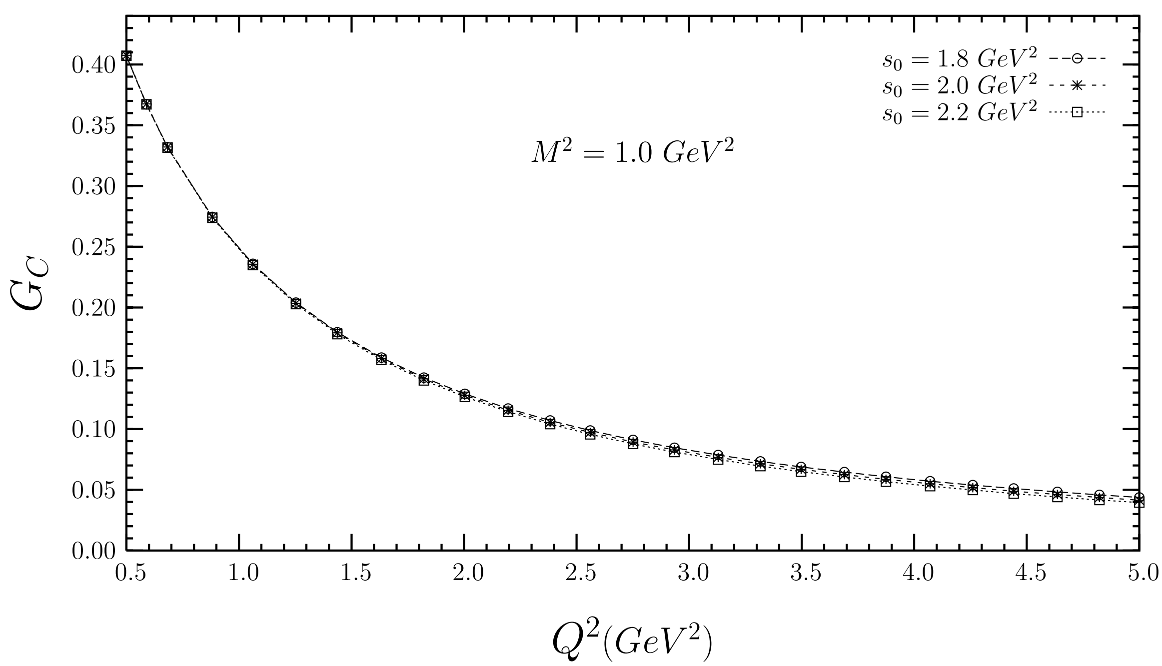

Figure 4:

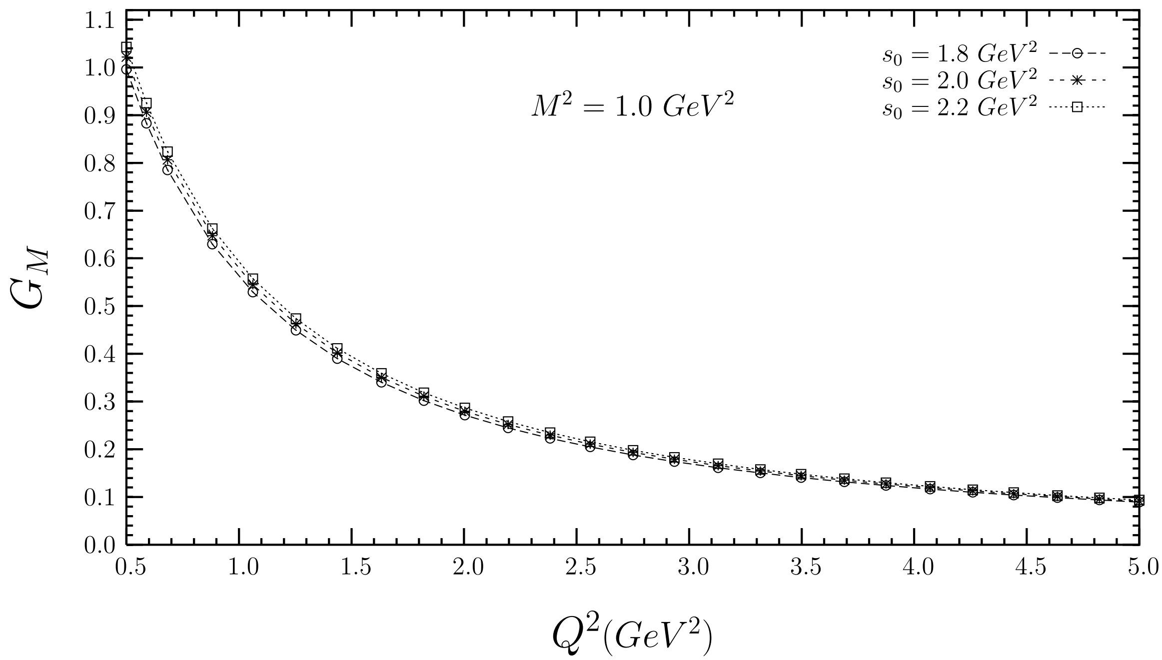

Figure 5:

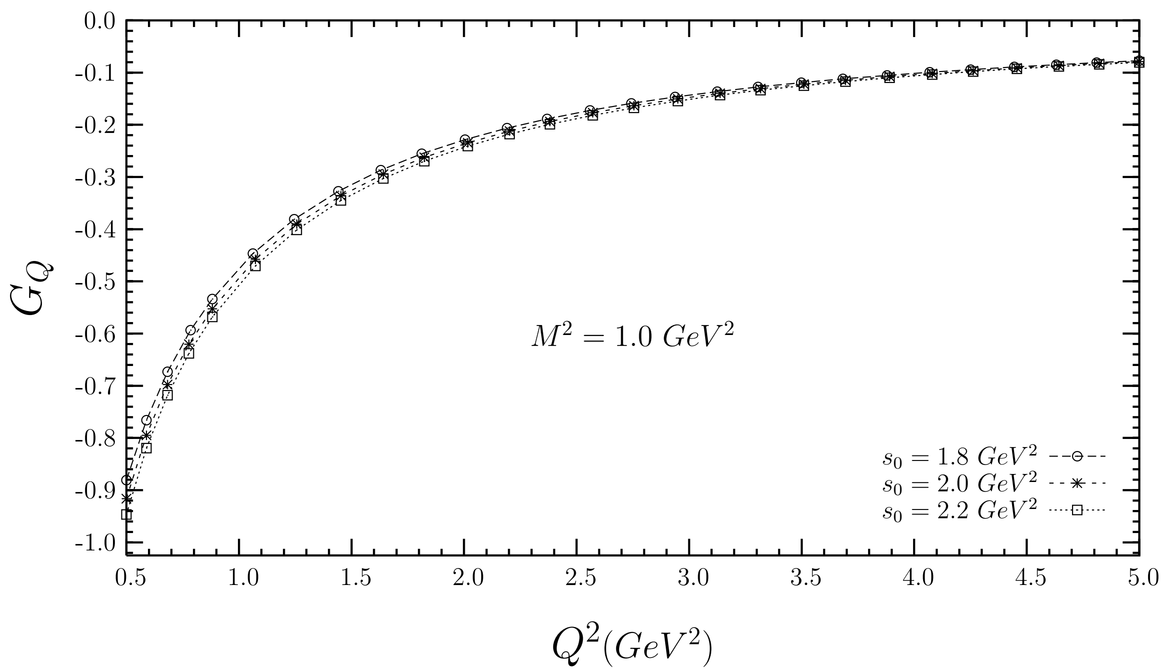

Figure 6: