Rotational covariance and light-front current matrix elements

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Light-front current matrix elements for elastic scattering from hadrons with spin 1 or greater must satisfy a nontrivial constraint associated with the requirement of rotational covariance for the current operator. Using a model $\rho$ meson as a prototype for hadronic quark models, this constraint and its implications are studied at both low and high momentum transfers. In the kinematic region appropriate for asymptotic QCD, helicity rules, together with the rotational covariance condition, yield an additional relation between the light-front current matrix elements.
I. INTRODUCTION

Light-front dynamics has found frequent application in particle and nuclear physics. First introduced by Dirac [3], it has the advantage that, of the ten generators of transformations for the Poincaré group, only three of them depend upon interaction dynamics. In particular, certain forms of Lorentz boosts are interaction independent. This is in contrast to the more traditionally used instant-form dynamics, which has four interacting generators, including all boosts.

However, as with all forms of relativistic dynamics, more generators than just the light-front “Hamiltonian” \( P^- = P^0 - \mathbf{P} \cdot \mathbf{n} \), where \( \mathbf{n} \) is a spatial unit vector, must depend nontrivially on the interaction, and these generators correspond to rotations about axes perpendicular to \( \mathbf{n} \). This is especially important when computing matrix elements of electromagnetic and weak currents in the front form. Because the current operator must have the transformation properties of a four-vector, and some of these transformations are interaction dependent, it must also depend upon the strong interaction. In particular, the current operator cannot satisfy all of the covariance requirements associated with transverse rotations without containing interaction dependent components. This is sometimes called the “angular condition” [1,2].

Rotational covariance has been studied in a variety of contexts. For hadrons composed of quarks, Terent’ev and others have examined the extent to which the current operator is uniquely determined without knowing its two-body components. For elastic electron-deuteron scattering, the lack of uniqueness modulo two-body currents has been explored by employing a variety of “schemes” to satisfy the rotational covariance requirement [3,7].

In this paper, we examine a specific rotational constraint for light-front current matrix elements of hadrons composed of quarks. The current operator must satisfy certain non-trivial commutation relations with the interacting generators of the Poincaré group. The requirement of current covariance for rotations about an axis perpendicular to \( \mathbf{n} \) gives rise to constraints on current matrix elements for elastic scattering from particles of spin 1 or
greater [54]. While form factors of the pion [10] and the nucleon [11] have frequently been
calculated, the particular test discussed here is not applicable because $j < 1$. We therefore
take a model of a $\rho$ meson as a prototype hadron to which to apply the rotational covariance
test.

II. ROTATIONAL COVARIANCE

In light-front dynamics, full rotational covariance implies a nontrivial set of conditions
which any hadronic model must in principle satisfy.

First, the state vector for the hadron must be an eigenstate of the total spin operator. This
condition is satisfied in light-front models which use a quantum mechanical Hilbert
space with a fixed number of particles [10,11]. Models using a field theory, in particular
those motivated by QCD, may not necessarily use rotationally covariant state vectors. At
high $Q^2$, this deficiency may be irrelevant if the corrections to rotational covariance fall as a
power of $Q^2$. At moderate and low $Q^2$, the issue may be important. In any event, we consider
here only models whose state vectors have the proper rotational covariance properties.

Second, the current operator $I^\mu(x)$ must satisfy the conditions of Lorentz covariance. If
$\Lambda^{\mu\nu}$ is the matrix for a homogeneous Lorentz transformation and $a^\mu$ is a spacetime transla-
tion, then

$$U(\Lambda, a)I^\mu(x)U(\Lambda, a)\dagger = (\Lambda^{-1})^\mu_\nu I^\nu(\Lambda x + a).$$

(1)

It must be conserved with respect to the four-momentum. If $P^\mu$ is the generator of spacetime
translations, then

$$g_{\mu\nu}[P^\mu, I^\nu(0)] = 0.$$

(2)

These constraints have two implications. First, it is possible to express the physical con-
tent of current matrix elements between any two states in terms of a limited number of
Lorentz invariant functions of the masses of the states and the square of the momentum
transfer. Second, the operator $I^\mu(0)$ must have in general a complicated structure, since it obeys nontrivial commutation relations with at least some generators which are interaction dependent.

To illustrate these two points, consider current matrix elements with spacelike momentum transfer. It has been shown \cite{7,9} that all spin matrix elements of the current four-vector can be computed from the set of matrix elements of $I^+(0)$ in a frame in which $q^+ = q^0 + \mathbf{q} \cdot \mathbf{n} = 0$. Alternatively, it means that all invariant form factors can be computed from matrix elements of $I^+(0)$. However, the covariance requirements imply that the matrix elements of $I^+(0)$ must satisfy a set of constraints. This is particularly relevant if a one-body operator has been used to compute the matrix elements. If a constraint involves transformations which use Poincaré generators which are non-interacting, it will in general be satisfied for matrix elements computed with one-body operators. However, constraints which involve transformations using interacting generators will in general not be satisfied with one-body current matrix elements.

An interaction-dependent constraint can be derived by requiring that Breit-frame matrix elements of the transverse current in a helicity basis vanish if the magnitude of the helicity flip is 2 or greater. For light-front current matrix elements $\langle \mathbf{p}'\alpha'\mid I^+(0)\mid \mathbf{p}\alpha \rangle$ corresponding to elastic electron scattering from a target of mass $M$ and spin $j$, the condition is expressed as follows:

$$\sum_{\lambda'\lambda} D_{\mu'\lambda'}^{(j)}(R_{ch}'(\mathbf{p}', M, \pi/2)) D_{\lambda\mu}^{(j)}(R_{ch}(\mathbf{p}, M, \pi/2)) = 0, \quad |\mu' - \mu| \geq 2. \quad (3)$$

In Eq. (3), $\mathbf{p} \equiv (p_\perp, p^+)$ is a light-front momentum, and the matrix element is evaluated in a frame where $q^+ = p'^+ - p^+ = 0$, and the perpendicular component of $\mathbf{p}'$ and $\mathbf{p}$ lies along the $x$ axis. The rotation

$$R_{ch} = R_{cf}(\mathbf{p}, M)R_y(\pi/2); \quad R'_{ch} = R_{cf}(\mathbf{p}', M)R_y(\pi/2), \quad (4)$$

where $R_{cf}$ is a Melosh rotation which, together with the rotation $R_y(\pi/2)$, transforms the state vectors from light-front spin to helicity. For inelastic excitation of a state with mass
$M'$ and spin $j'$, Eq. (3) is modified only by the use of $M'$ and $j'$ in the rotation matrices for
the final state. For elastic scattering, Eq. (3) is applicable only to states with $j \geq 1$. For a
spin-1 particle, Eq. (3) can be expressed explicitly in terms of individual light-front matrix
elements as follows [5]:

$$
\Delta(Q^2) \equiv (1 + 2\eta)I_{1,1} + I_{-1,-1} - \sqrt{8\eta}I_{1,0} - I_{0,0} = 0,
$$

(5)

where $I_{\mu',\mu} \equiv \langle \tilde{p}'\mu'|I^+(0)|\tilde{p}\mu \rangle$ is the matrix element of $I^+(0)$:

$$
\begin{align*}
\mathbf{p}'_\perp &= -\mathbf{p}_\perp = \frac{1}{2}\mathbf{q}; & \quad p'^+ &= p^+ = \sqrt{M^2 + \frac{1}{4}q^2},
\end{align*}
$$

(6)

$Q^2 = -q^2$ is the square of the four-momentum transfer, $\eta \equiv Q^2/4M^2$.

Earlier studies of the deuteron form factors using models with one-body currents and
rotationally covariant state vectors indicate that $\Delta(Q^2)$ can be relatively small for low and
moderate $Q^2$, though, not surprisingly, $\Delta(Q^2)$ is an increasing function of $Q^2$ [5,7]. An
important feature of the deuteron is that the characteristic nucleon momentum is very small
compared to a nucleon mass.

For most quark models of hadrons, the characteristic constituent momentum is not small
with respect to the quark mass, and questions of rotational covariance therefore require a
separate investigation. There have been previous studies of the pion [10] and the nucleon [11]
form factors using models with one-body currents and rotationally covariant state vectors.
Equation (5) has no counterpart for spin zero and spin $\frac{1}{2}$, so the dynamical aspect of ro-
tational covariance was not addressed in those works. Some other conclusions from those
studies which are relevant to the discussion below include

1. For small quark masses (10 MeV), relativistic effects such as those of Melosh rotations
can be substantial, even at low $Q^2$.

2. In the limit $m_q \to 0$, it can be shown [12] that $Q^2F_{\pi}(Q^2) \to \text{const}$ as $Q^2 \to \infty$.

In what follows, we examine a model for a $\rho$ meson similar to those used for the pion
published earlier, in light of the rotational covariance condition (3) at both low and high
$Q^2$. 

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III. THE MODEL

The model $\rho$ meson is composed of a valence quark and an antiquark. Since it must be color-antisymmetric and flavor-symmetric (isospin 1), the space-spin wave function must be symmetric. For the ground state, we take the coupled states $^3S_1 - ^3D_1$. This is the same space-spin combination as that of a deuteron composed of two nucleons. Since the details of such a deuteron form factor calculation are discussed extensively elsewhere [7], only the unique features of the quark model will be presented here. The mass operator is

$$M^2 = M_0^2 + 2mU + U_0,$$

where $M_0^2 = 4(m^2 + k^2)$ is the non-interacting mass operator, $m$ is the quark mass and $k$ is the relative three-momentum. The potential $U$ is a harmonic-oscillator potential:

$$U = \frac{1}{2}\kappa\mathbf{r}^2, \quad \mathbf{r} \equiv i \nabla_k,$$

and $U_0$ is a constant. The $S$- and $D$-state wave functions are Gaussians, just as for a non-relativistic harmonic-oscillator problem:

$$\phi_0(k) = N_S e^{-k^2/2\alpha^2};$$
$$\phi_2(k) = N_D k^2 e^{-k^2/2\alpha^2}.$$  \(9\)

The interacting mass eigenvalue is

$$M^2 = 4(m^2 + 3\alpha^2) + U_0.$$  \(10\)

In an extensive study of mesons using a nonrelativistic quark model, Godfrey and Isgur obtained $D$-state admixtures with amplitude 0.04 for an excited $\rho$-meson state, with no reported admixture for the $\rho$ (750) [13]. To test sensitivity, we use a $D$-state admixture amplitude coefficient of 0.04; this is an extreme choice, but in the end, the results differ little from those obtained by ignoring the $D$-state admixture entirely. An oscillator parameter value $\alpha = 0.45$ GeV/$c$ has been extracted from the Godfrey-Isgur results [14].
IV. FORM FACTORS

Matrix elements of matrix elements of $I^+(0)$ can be written as \[7\]

$$
\langle \mathbf{p}' \mu' | I^+(0) | \mathbf{p} \mu \rangle = \langle p_1' + \frac{1}{2} q_\perp \mu' | I_{\text{quark}}(0) | p_1' - \frac{1}{2} q_\perp \mu_1 \rangle \int \frac{d \mathbf{k}}{(2\pi)^3} \left| \frac{\partial (\mathbf{p} \mathbf{k})}{\partial (\mathbf{p}_1 \mathbf{p}_2)} \right|^\frac{1}{2} \left| \frac{\partial (\mathbf{p}'_1 \mathbf{p}'_2)}{\partial (\mathbf{p}' \mathbf{k}')} \right|^\frac{1}{2} \times D^{(\frac{3}{2})}_{\mu_1 \mu_1} [R_{\text{cf}}(\mathbf{k}')] D^{(\frac{3}{2})}_{\mu_2 \mu_2} [R_{\text{cf}}(-\mathbf{k}')] \\
\times \langle \frac{1}{2} \mu'_1 \frac{1}{2} \mu'_2 | 1 \mu'_3 \rangle \langle l \mu'_1 1 \mu'_3 | 1 \mu' \rangle Y_{l \mu'}(\hat{\mathbf{k}}') \phi_{l'}(k') \\
\times D^{(\frac{3}{2})\dagger}_{\mu_1 \mu_1} [R_{\text{cf}}(\mathbf{k})] D^{(\frac{3}{2})\dagger}_{\mu_2 \mu_2} [R_{\text{cf}}(\mathbf{k})] \\
\times \langle \frac{1}{2} \mu_1 \frac{1}{2} \mu_2 | l \mu_3 \rangle \langle l \mu_1 \mu_2 | l \mu \rangle Y_{l \mu}(\hat{\mathbf{k}}) \phi_l(k). \tag{11}
$$

The internal kinematics in the integral are

$$
p_1' + = p_1^+ = \sqrt{m^2 + \frac{1}{4} q_\perp^2}; \quad p_2' + = p_2^+; \quad p_1' \perp = p_1' + q_\perp; \quad p_2' \perp = p_2' + q_\perp \\
\mathbf{k}_\perp = (1 - \xi) p_1 \perp - \xi q_\perp; \quad \mathbf{k}_1 = \mathbf{k}_\perp + (1 - \xi) q_\perp; \quad \xi \equiv p_1^+ / (p_1^+ + p_2^+) \\
k_3 = (\xi - \frac{1}{2}) \sqrt{\frac{m^2 + k_\perp^2}{\xi(1 - \xi)}}; \quad k_3' = (\xi - \frac{1}{2}) \sqrt{\frac{m^2 + k_\perp'^2}{\xi(1 - \xi)}}; \\
\mathbf{k} = (\mathbf{k}_\perp, k_3); \quad \mathbf{k}' = (\mathbf{k}_\perp, k_3'). \tag{12}
$$

The three elastic form factors $G_1$, $G_2$ and $G_3$ can be expressed in terms of the matrix elements $I_{\mu', \mu}$ as follows \[7\]:

$$
G_0(Q^2) = \frac{1}{2(1 + \eta)} \left[ (1 - \frac{2}{3} \eta)(I_{1,1} + I_{0,0}) + \frac{5}{3} \sqrt{8 \eta I_{1,0}} - \frac{1}{3}(1 - 4 \eta) I_{1,-1} \right] \\
G_1(Q^2) = \frac{1}{(1 + \eta)} \left[ I_{1,1} + I_{0,0} - I_{1,-1} - (1 - \eta) \sqrt{\frac{2}{\eta} I_{1,0}} \right] \\
G_2(Q^2) = \frac{\sqrt{8}}{3(1 + \eta)} \left[ -\frac{\eta}{2}(I_{1,1} + I_{0,0}) + \sqrt{2 \eta I_{1,0}} - (1 + \frac{1}{2} \eta) I_{1,-1} \right]. \tag{13}
$$

The right-hand sides in Eq. \[13\] are not unique. One can always replace one of the $I_{\mu', \mu}$, or linear combinations of them, with a combination which satisfies the rotational covariance condition \[5\]. A common procedure has been to choose a particular combination of $I_{\mu', \mu}$ as calculated from one-body current operators, and eliminating the remaining terms from
Eq. (13) via the rotational covariance condition. By implication, the eliminated terms depend upon two-body current operators. Thus, for different choices of one-body matrix elements, each form factors $G_{i}(Q^{2})$ will differ by a multiple of $\Delta(Q^{2})_{\text{one-body}}$, which is never zero.

V. Low $Q^{2}$

The requirement of rotational covariance can be studied at low momentum transfers by examining the behavior of the magnetic and quadrupole moments $\mu$ and $\bar{Q}$ and the charge radius. They are related to the form factors $G_{i}(Q^{2})$ as follows:

$$
\mu \equiv \lim_{Q^2 \to 0} G_{1}(Q^{2})
$$

$$
\bar{Q} \equiv \lim_{Q^2 \to 0} 3\sqrt{2} \frac{G_{2}(Q^{2})}{Q^{2}}
$$

$$
\langle r^{2} \rangle \equiv \lim_{Q^2 \to 0} \frac{6}{Q^{2}} [1 - G_{0}(Q^{2})].
$$

Extracted values of $\langle r^2 \rangle^{\frac{1}{2}}$ are shown in Table I for quark masses of 10, 300 and 1000 MeV. Also shown is the effect of including or leaving out the Melosh rotations, which gives a measure of the size of relativistic effects. In addition, the quantity

$$
\delta \equiv \lim_{Q^2 \to 0} \frac{\Delta(Q^{2})}{1 - G_{0}(Q^{2})}
$$

gives a measure of the sensitivity to rotational covariance uncertainty. Already one can see from this table corresponding to very low $Q^2$ that simply raising the value of the quark mass is not the same as the nonrelativistic limit. That limit depends upon the quark mass, the value of $Q^2$, the momentum scale $\alpha$, and the composite mass. For comparison, we also show results using the same parameter, except that the $\rho$ meson is given a fictitious value of 2 GeV. In this last case, especially for a quark mass of 1 GeV, one can see that both the Melosh rotations (the measure of relativistic effects) and the rotational covariance parameter $\delta$ are small.
VI. MODERATE $Q^2$

In Figs. 1, 2 and 3 we show the calculated results for $G_0(Q^2)$, as obtained using Eq. (13). The relativistic effects, as characterized by turning the internal Melosh rotations on and off, are largest for the smallest quark mass. For all three choices of quark mass, the rotational covariance uncertainty function $\Delta(Q^2)$ becomes comparable to $G_0$, and hence the current matrix elements themselves, in the region 1–2 GeV/$c^2$.

In Figs. 4, 5 and 6 we show calculated results for $G_0(Q^2)$ using a fictitious $\rho$-meson mass of 2 GeV. In all three cases, the relativistic effects are smaller than the corresponding cases where $M_\rho = 750$ MeV, but the nonrelativistic limit is still not really reached until the quark mass is considerably larger than the momentum scale $\alpha$. For this choice of meson mass, and for all three choices of quark mass, the covariance function $\Delta(Q^2)$ is much smaller for the same range of $Q^2$ than in the previous three figures.

From the results shown here, along with those of other parameter sets, it becomes clear that the rotational covariance uncertainty function $\Delta(Q^2)$ becomes comparable to $G_0$, and hence the current matrix elements themselves, when $\eta = Q^2/4M^2$ is of order unity. Thus, for a $\rho$ meson with physical mass 750 MeV, the breakdown of rotational covariance occurs in the region 1–2 GeV/$c^2$. This can be understood from the fact that the dimensionless argument of the Melosh rotations in Eq. (3) is $Q/2M$, which manifests itself in terms of the $\eta$ factors in Eq. (5). The dynamical nature of the rotational covariance condition is contained in the presence of the interacting mass $M$.

Note also that, for elastic scattering, the current matrix elements $I_{\mu',\mu}$ depend upon the quark mass $m$ and the momentum transfer, but they do not depend upon the composite mass $M$. The internal Melosh rotations, which give a measure of size of relativistic effects, depend upon the quark mass, but the composite mass enters only at the point of computing form factors and evaluating the rotational covariance condition.
VII. ASYMPTOTIC BEHAVIOR

As noted above, it has been shown that, for the pion form factor, models such as the one used here have the property that $Q^2 F_\pi(Q^2) \to \text{const}$ as $Q^2 \to \infty$ if $m_q = 0$. For a model $\rho$ meson, the differing feature is the presence of an overall spin index and some momentum-independent Clebsch-Gordan coefficients. We therefore expect that, as for the case of the pion form factor \cite{12}, $Q^2 I_{\mu',\mu} \to \text{const}$ (perhaps dependent upon $\mu', \mu$) as $Q^2 \to \infty$ for $m_q = 0$ and any $\mu', \mu$. On the other hand, power counting rules of perturbative QCD \cite{16} predict that the matrix element $I_{0,0}$ dominates as $Q^2 \to \infty$, and that $I_{1,0}$ is suppressed by one power of $Q$ and $I_{1,-1}$ by two powers of $Q$. Thus, a constituent quark model with one-body currents only cannot reproduce the asymptotic limit.

While simple models may fail to describe the asymptotic limit appropriate for perturbative QCD, we note that the rotational covariance condition takes a simple form at very high $Q^2$. In the $Q^2 \to \infty$ limit of Eq. (5), $I_{1,-1}$ drops out due to power suppression. The remaining terms give

$$2\eta I_{1,1} - \sqrt{8}\eta I_{1,0} - I_{0,0} = 0,$$

The Breit frame $(1,1)$ matrix element of the transverse current in a helicity basis is identically zero \cite{16}. The light-front matrix element $I_{1,1}$ is not identically zero, but is suppressed by two powers of $Q$ relative to a specific combination of $I_{0,0}$ and $Q I_{1,0}$.

Note also that, in Eq. (16), the factors of $\eta$ which correspond to the dynamical nature of the rotational covariance condition are now very large. At the same time, the gluon-exchange terms used to derive the power-counting helicity rules in perturbative QCD correspond to two-body currents in a constituent model such as that presented here. At high $Q^2$, therefore, the $\eta$ factors, the dynamics of perturbative QCD and rotational covariance are linked in a way which cannot be described in a constituent model with one-body currents.
VIII. CONCLUSIONS

The requirement of rotational covariance for matrix elements of electromagnetic currents is nontrivial to satisfy for elastic scattering from systems with spin $j \geq 1$. For the case of the deuteron, the fact that its structure is essentially nonrelativistic (all masses are large compared to the momentum scale of the system) suggests that the rotational covariance requirement can be satisfied in a satisfactory quantitative way using only one-body current matrix elements. For hadrons composed of quarks, typical quark masses are not small compared to typical momentum scales, and issue of rotational covariance therefore must be studied separately. In this paper, we have investigated the behavior of current matrix elements in a simple model of a $\rho$ meson. Our results using variable input parameters indicate a breakdown of rotational covariance of current matrix elements when $\eta = Q^2/4M^2$ is of order unity. The dynamical nature of the rotational covariance constraint is reflected in the presence of the interacting mass eigenvalue of the composite particle. In addition, rotational covariance implies a specific power-law relation among the spin matrix elements at high $Q^2$. That relation is consistent with the power-counting rules of perturbative QCD, which in turn are derived from gluon-exchange contributions that correspond to two-body currents in a constituent-quark framework. The quark model used in this paper does not contain such two-body currents, and also does not satisfy the high-$Q^2$ power-law relation.

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TABLE I. Calculated r.m.s. charge radius, together with the dimensionless rotational covariance parameter $\delta$, for various model parameters.

| $M_\rho$ (MeV) | $m_q$ (MeV) | Melosh | $\langle r^2 \rangle^{1/2}$ (fm) | $\delta$ |
|----------------|-------------|--------|-------------------------------|---------|
| 750            | 10          | ON     | 1.17                          | .03     |
|                |             | OFF    | 0.76                          | .16     |
| 300            | 10          | ON     | 0.61                          | .17     |
|                |             | OFF    | 0.54                          | .36     |
| 1000           | 10          | ON     | 0.51                          | .36     |
|                |             | OFF    | 0.50                          | .42     |
| 2000           | 10          | ON     | 1.11                          | .11     |
|                |             | OFF    | 0.66                          | .05     |
| 300            | 10          | ON     | 0.45                          | .19     |
|                |             | OFF    | 0.37                          | .11     |
| 1000           | 10          | ON     | 0.33                          | .04     |
|                |             | OFF    | 0.32                          | .15     |
FIGURES

FIG. 1. Composite form factor $G_0(Q^2)$ computed using a quark mass of 10 MeV and a $\rho$-meson mass of 750 MeV. The solid curve denotes the full calculation and the dashed curve corresponds to the calculation where the Melosh rotations are turned off. The dot-dashed curve describes the covariance function $\Delta(Q^2)$ defined in Eq. (5).

FIG. 2. Composite form factor $G_0(Q^2)$ computed using a quark mass of 300 MeV and a $\rho$-meson mass of 750 MeV. The legend is the same as that of Fig. 1.

FIG. 3. Composite form factor $G_0(Q^2)$ computed using a quark mass of 1 GeV and a $\rho$-meson mass of 750 MeV. The legend is the same as that of Fig. 1.

FIG. 4. Composite form factor $G_0(Q^2)$ computed using a quark mass of 10 MeV and a $\rho$-meson mass of 2 GeV. The legend is the same as that of Fig. 1.

FIG. 5. Composite form factor $G_0(Q^2)$ computed using a quark mass of 300 MeV and a $\rho$-meson mass of 2 GeV. The legend is the same as that of Fig. 1.

FIG. 6. Composite form factor $G_0(Q^2)$ computed using a quark mass of 1 GeV and a $\rho$-meson mass of 2 GeV. The legend is the same as that of Fig. 1.
$m_q = 10 \text{ MeV}; \ M_\rho = 750 \text{ MeV}$
$m_q = 300$ MeV; $M_\rho = 750$ MeV
$m_q = 1000 \text{ MeV}; M_\rho = 750 \text{ MeV}$
$m_q = 10 \text{ MeV}; \ M_\rho = 2000 \text{ MeV}$
$m_q = 300$ MeV; $M_\rho = 2000$ MeV
\( m_q = 1000 \text{ MeV}; \quad M_\rho = 2000 \text{ MeV} \)