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Pion transition form factor in the Regge approach

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We explore the BaBar puzzle within the Regge approach. After reviewing the chiral quark models in applications to PDF and PDA of the pion, we argue that variants of these models, fulfilling the chiral anomaly, may in fact violate the second Terazawa-West unitarity bound, which is based on unverified assumptions for the real part of the amplitude. Consequently, the transition form factor need not vanish at large values of the photon virtuality. Then we show that the experimental data may be properly explained with incomplete vector-meson dominance in a simple model with one state, as well as in more sophisticated radial Regge models including infinitely many states. We also consider the experimental constraint from the rare $Z \to \pi_0 \gamma$ decay, which is comfortably satisfied in our approach. Finally, we point out that the photon momentum asymmetry parameter may noticeably influence the precision fits to the data.

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1. Introduction

After the release of the BaBar data [1] for the pion-photon transition form factor, $F_{\pi^0\gamma\gamma^*}$, our community is in deep shock, as the conventional approach to the gold-plated exclusive process at very high Euclidean momenta Q, based on 1. factorization and 2. (leading-twist) pQCD evolution [2 – 12], seems to be invalid [13, 14]. Indeed, the quantity $Q^2 F_{\pi^0\gamma\gamma^*}(Q^2)$ goes visibly above the famous Brodsky-Lepage limit,

$$Q^2 F_{\pi^0\gamma\gamma^*}(Q^2) \to \frac{2f_\pi}{N_c}\int_0^1 dx \frac{\phi_{\rm as}(x)}{x} = \frac{6f_\pi}{N_c} = 2f_\pi, \tag{1.1}$$

at momenta $Q^2 > 15$ GeV^2.

Several ideas, abandoning assumptions 1. and 2., have been proposed to solve the "BaBar problem". Radyushkin [15] and Polyakov [16] advocated that the possible nonvanishing of the PDA at the end points (as found by the present authors in chiral quark models at the low-energy quark-model scale [17]), together with essentially switched-off evolution and regulated quark propagators, is capable of reproducing the data in the CLEO and BaBar domain. In this approach $F_{\pi^0\gamma\gamma^*} \sim \log(Q^2/\mu^2)/Q^2$, with the log indicating the breaking of factorization (note that the same asymptotics follows in the Spectral Quark Model (SQM), cf. Eq. (14.1) of [18]). Dorokhov [19 – 21] proposed the use of the fixed-mass chiral quark model to evaluate the triangle diagram of Fig. 1, which is capable of reproducing the data at high Q^2, however the fitted value of the constituent quark mass is very low, $M \sim 135$ MeV. In this the form $F_{\pi^0\gamma\gamma^*} \sim [\log(Q^2/\mu^2)]^2/Q^2$. A possible need of higher-twist terms has been brought up in [22]. The calculation of [23] in the nonlocal chiral quark model inspired by the instanton-liquid model of QCD produced the result in agreement with the data at lower values of Q^2 and complying to the limit (1.1). On the other hand, Dorokhov [24] considered a modified pion-meson vertex [25] in the nonlocal model and found agreement in the whole available data range, albeit again with a very low constituent quark mass. Asymptotically, in this approach $F_{\pi^0\gamma\gamma^*} \sim \log(Q^2/\mu^2)/Q^2$. The influence of nonperturbative gluonic components of the pion has been considered in [26]. Concerns on the validity of factorization were discussed in [27]. It was also shown that the pronounced growth of $Q^2 F_{\pi^0\gamma\gamma^*}$ between 10 and 40 GeV^2 cannot be explained via higher-order pQCD corrections at the NNLO level [12]. Further developments may be found in [28 – 31]

This talk is based on [32].

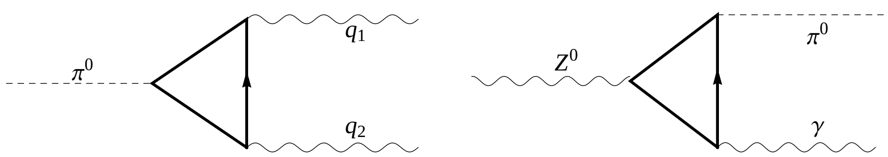

Figure 1: The quark-model diagrams used to evaluate the pion transition form factor (left) and the $Z_0 \to \pi^0\gamma$ decay (right). The crossed diagrams not shown.

2. General constraints

Consider the process of the left part of Fig. 1 with the general kinematics

$$q_1^2 = -\frac{1+A}{2}Q^2,\ q_2^2 = -\frac{1-A}{2}Q^2,\ \ -1 \le A \le 1. \tag{2.1}$$

The chiral anomaly [33, 34] fixes $F_{\pi^0\gamma^*\gamma^*}(Q^2=0,A) = \frac{1}{4\pi^2 f_\pi}$.

On the other, the high-Q^2 behavior of $F_{\pi^0\gamma\gamma^*}$ is formally limited by the Terazawa-West (TW) [35 – 37] unitarity bounds, recently brought up by Dorokhov [21] and further elaborated in [32]. The derivation of the TW bounds uses the Schwarz inequality involving sums of the matrix elements $\langle 0|J_\mu(0)|n\rangle$ and $\langle \pi^a(q)|J_\mu(0)|n\rangle$, entering the vacuum polarization and the parton distribution function (PDF) of the pion. Schematically, one has $|\langle\pi|JJ|0\rangle| \le (|\langle 0|JJ|0\rangle||\langle\pi|JJ|\pi\rangle|)^{1/2}$. Then one derives the first bound,

$$\mathrm{Im} F_{\pi^0\gamma\gamma^*}(q^2) = \mathcal{O}(1/\sqrt{q^2}) \quad \text{(TW I)} \tag{2.2}$$

holding at *time-like* momenta, $q^2 > 4m_\pi^2$. If there are *no polynomial terms* in the real part of $F_{\pi^0\gamma\gamma^*}$ (which is an *assumption*), then $\left|F_{\pi^0\gamma\gamma^*}(q^2)\right| = \mathcal{O}(1/\sqrt{q^2})$. A dispersion relation yields [37] the second bound,

$$\left|F_{\pi^0\gamma\gamma^*}(Q^2)\right| = \mathcal{O}(1/Q) \quad \text{(TW II)} \tag{2.3}$$

valid for all momenta, also large *space-like* momenta Q. The constant in the bound may be explicitly given [37] in terms of the photon spectral density and the pion structure function,

$$\left|F_{\pi^0\gamma\gamma^*}(Q^2)\right| < \frac{2\sqrt{\Pi(\infty)}}{Q}\int_0^1 dx \sqrt{\frac{F_1(x,Q^2)}{x(1-x)}} \tag{2.4}$$

where $\Pi(s) = s/(16\pi^3\alpha_{\mathrm{QED}}^2)\sigma_{e^+e^-\to\mathrm{hadrons}}(s)$ and $\Pi(\infty) = 1/(12\pi^2)\sum_i e_i^2$, with e_i denting the quark charges. With the SMRS [38] and GRV [39] parameterizations for F_1 we obtain (for Q^2 in the range $10-40\ \mathrm{GeV}^2$)

$$|F_{\pi^0\gamma\gamma^*}(Q^2)| < \frac{0.85(1)}{Q} \quad \text{(LO)}, \quad < \frac{0.75(1)}{Q} \quad \text{(NLO)}, \tag{2.5}$$

where LO and NLO refer to Leading Order and Next-to-Leading Order in the QCD evolution, respectively. Thus the TW II bound is “inefficient” in the present experimental range, as it goes an order of magnitude above the BaBar data.

Finally, another interesting experimental bound comes from the rare $Z \to \pi^0\gamma$ decay [40], which probes the time-like value $q^2 = M_Z^2$. The process is shown in the right part of Fig. 1. Since only the vector coupling of the Z^0 boson to the quark contributes,

$$\frac{F_{Z\to\pi^0\gamma}(q^2)}{F_{Z\to\pi^0\gamma}(0)} = \frac{F_{\pi^0\gamma^*\gamma}(q^2)}{F_{\pi_0\gamma^*\gamma}(0)}. \tag{2.6}$$

The experimental limit given by the Particle Data Group [41], $\Gamma(Z^0\to\pi^0\gamma) < 5\times10^{-5}\Gamma_{\mathrm{tot}}(Z^0) = 10.25\times10^{-5}\mathrm{GeV}$, implies

$$|F_{Z^0\to\pi^0\gamma}(M_Z^2)/F_{Z^0\to\pi^0\gamma}(0)| < 0.17. \tag{2.7}$$

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3. Subtracted dispersion relation and violation of TW II

The assumption of the absence of the polynomial terms, necessary for TW II to hold, is equivalent to the validity of the unsubtracted dispersion relation for $F_{\pi^0\gamma^*\gamma}$. Clearly, pQCD with factorization leads to $F_{\pi^0\gamma^*\gamma}$ vanishing at $Q \to \infty$, however, since now these assumptions are questioned (see the Introduction), one may consider the situation where subtraction constants are necessary in the dispersion relation for $F_{\pi^0\gamma^*\gamma}$ [32]. Note that even if the form factor vanishes at infinity, one can write a subtracted relation

$$F(t) - F(0) = \frac{1}{\pi}\int_{s_0}^{4\Lambda^2} \frac{t}{s}\frac{\mathrm{Im}F(s)}{s-t}ds + \frac{1}{\pi}\int_{4\Lambda^2}^{\infty} \frac{t}{s}\frac{\mathrm{Im}F(s)}{s-t}ds. \tag{3.1}$$

If Λ is large ($\Lambda^2 > Q^2$), the second term is very slowly varying with Q^2 and mimics a constant. In particular, for $\mathrm{Im}F(s) \sim 1/\sqrt{s}$ it behaves as $1/\Lambda + \mathscr{O}(1/Q)$. Thus the appearance of the constant term need not be taken as a fundamental problem, as it may represent the unknown high-energy data. We will analyze below a quantitative lower bound for a possible high energy mass scale.

The bottom line is that one may well consider the case where $F_{\pi^0\gamma^*\gamma}(Q^2)$ does not vanish at $Q^2 \to \infty$. Below we will provide an explicit field-theoretic example where this is the case (the Georgi-Manohar (GM) model [42]). Motivated by this we will also incorporate a constant term in the fits of the BaBar data.

4. Mini-review of chiral quark models

Before discussing the GM model let us briefly review the chiral quark models and, in particular, their predictions for the *soft* matrix elements entering the (factorized) high-energy processes. The purpose is to convince the reader that these models, based on *chiral symmetry breaking* as the key dynamical ingredient, are very successful for numerous processes involving the pions and photons. Chiral quark models are cast in a covariant Lagrangian form and in general exhibit no factorization. They carry relatively few parameters, traded for f_π, m_π, ... The large-N_c limit is implicit, however, confinement is absent and one needs to be careful not to open the $q\bar{q}$ production threshold. This limits the applicability range, however, no problems arise with space-like external momenta, as in the present analysis of the BaBar puzzle. Finally, predictions of chiral quark models hold at a low-energy *quark-model* scale [43] and QCD evolution is necessary to reach higher scales of experiments or lattice simulations.

We start with the valence PDF of the pion, where the NJL model gives [44]

$$q(x) = 1 \tag{4.1}$$

(note that the same *constant* PDF for two constituents follows from AdS/CFT approach of [45]). Arguments based on the momentum sum rule [43] determine the renormalization scale where Eq. (4.1) holds: $\mu_0 \sim 320$ MeV. Then at the scale $\mu = 2$ GeV the valence quarks carry 47% of the momentum of the pion, as requested by experimentally motivated parameterizations. The value of the coupling constant at the quark-model scale is $\alpha(\mu_0)/\pi = 0.68$. After the evolution the NJL prediction agrees remarkably well with the experimental data (left panel of Fig. 2).

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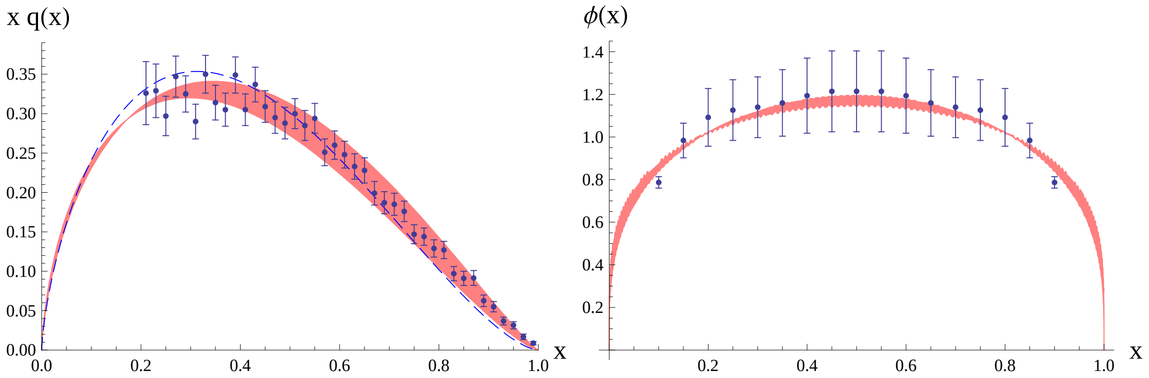

Figure 2: Left: The NJL prediction for the valence PDF of the pion evolved to the experimental scale $\mu = 4$ GeV (band). The data points come from the analysis of the Drell-Yan data from the E615 experiment [46]. The dashed line shows the reanalysis of the data from [47]. Right: The NJL prediction for the PDA of the pion evolved to the lattice scale $\mu = 0.5$GeV (band) compared to the transverse lattice data [48]. The width of the bands indicates the uncertainty in the quark-model scale μ_0.

In the NJL model, the PDA of the pion is also constant at the quark-model scale [17],

$$\phi(x) = 1 \tag{4.2}$$

(note that this result is different from the AdS/CFT prediction of [45], $\phi(x) \sim \sqrt{x(1-x)}$, it is also far from the asymptotic form $6x(1-x)$). The PDA (4.2) does not vanish at the end points, which is the focal point of [15, 16]. However, the LO ERBL evolution makes $\phi(x)$ vanish at the end-points, with the form $\phi(x) \sim x^{2C_F/\beta_0 \log[\alpha(\mu_0)/\alpha(\mu)]}$ near $x = 0$ [2, 43]. Thus maintaining the nonvanishing at the end points requires switching off the QCD evolution [15].

Finally, we mention the NJL results for the gravitational form factor of the pion, discussed in detail in [49]. The quark-model relation holds for the radii related to the spin-2 gravitational form factor and the charge form factor, $\langle r^2\rangle_\Theta = \frac{1}{2}\langle r^2\rangle_V$. Therefore matter is more concentrated than charge, in agreement with the recent AdS/CFT results [50]. Note that there is no contradiction between having a constant vertex function and a finite pion size. This issue is tightly linked with the fulfillment of the electromagnetic and chiral Ward identities, as discussed, e.g., in [51].

5. Violating TW II

We now recall a model which reproduces the results of the previous section, but violates TW II: the Georgi-Manohar (GM) model [42]. It is obtained from the NJL model by carrying out the chiral rotation (which is innocuous) and then introducing g_A of the quark which may be different from unity [52]. The Lagrangian of the GM model is

$$L = \bar{q}\left(i\partial\!\!\!/ + g_A^Q A\!\!\!/\gamma_5 - M\right)q + \frac{f^2}{4}\mathrm{Tr}\left(\partial_\mu U^\dagger \partial^\mu U\right) + \mathrm{WZW}, \tag{5.1}$$

where WZW stands for the Wess-Zumino-Witten term, while

$$A_\mu = \frac{i}{2}(u^\dagger \partial_\mu u - u\partial_\mu u^\dagger), \;\; u = e^{i\vec{\pi}\cdot\vec{\tau}/(2f)}, \;\; U = u^2. \tag{5.2}$$

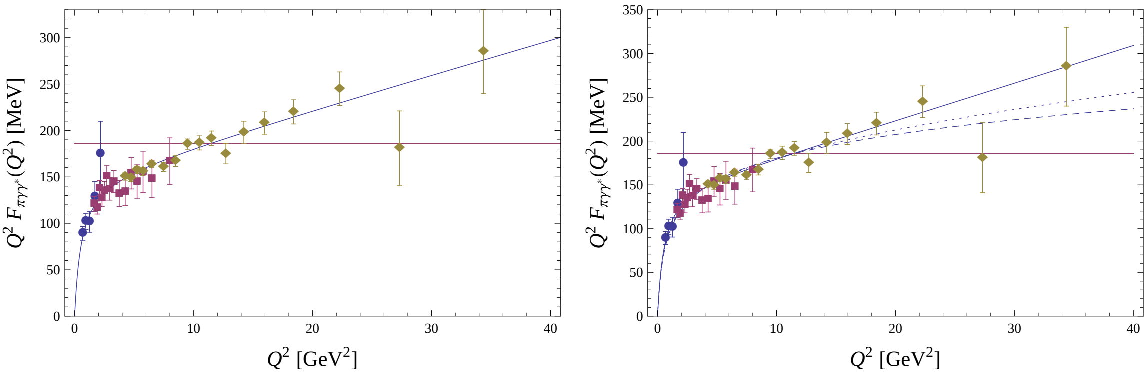

Figure 3: Left: predictions of the one-state IVMD model (line). Right: predictions of various Regge models: the Veneziano-Dominguez model $b = 1.81$ (dashed line), model with the first pole separated and fixed $b = 1.5$ (dotted line), and the subtracted Regge model (solid line). The dots, squares, and diamonds correspond to the CELLO [55], CLEO [56], and BaBar [1] data, correspondingly.

The resulting pion-photon transition form factor is

$$F_{\pi^0\gamma\gamma^*}(Q^2) = \frac{1}{4\pi^2 f_\pi} + \frac{g_A^Q}{4\pi^2 f_\pi}\left[G(Q^2)-1\right], \quad G(Q^2) = \frac{2M^2}{Q^2}\int_0^1 \frac{dx}{x}\log\left[1+x(1-x)\frac{Q^2}{M^2}\right] \tag{5.3}$$

As we can see, the anomaly is satisfied, but for $g_A^Q \neq 1$ *no vanishing* at $Q^2 \to \infty$ occurs,

$$F_{\pi^0\gamma\gamma^*}(Q^2) = \frac{1-g_A^Q}{4\pi^2 f_\pi} + \frac{g_A^Q M^2}{4\pi^2 f_\pi}\frac{\left[\log(Q^2/M^2)\right]^2}{Q^2} + \ldots \tag{5.4}$$

In the Spectral Quark Model [18] one has instead

$$G(Q^2) = \frac{1}{3}\left[\frac{2m_\rho^2}{m_\rho^2+Q^2} + \frac{m_\rho^2}{Q^2}\log\frac{m_\rho^2+Q^2}{m_\rho}\right], \quad F_{\pi^0\gamma\gamma^*} = \frac{1-g_A^Q}{4\pi^2 f_\pi} + \frac{g_A^Q m_\rho^2}{12\pi^2 f_\pi}\frac{\log\frac{Q^2}{m_\rho^2}}{Q^2} + \ldots \tag{5.5}$$

With $g_A^Q = 1$ this model fulfills the result of [15] with the mass scale m_ρ.

Unfortunately, precise chiral-quark-model fits of $F_{\pi^0\gamma\gamma^*}$ based on above formulas in the *whole* Q^2 range are not satisfactory, so in the following we proceed to more general analyses based on Incomplete Vector Meson Dominance (IVMD) [53, 54].

6. Incomplete VMD and Regge models

As shown in the previous Section, it is possible to formulate a field-theoretic model consistent with all formal requirements which violates TW II, leading to an asymptotically constant $F_{\pi^0\gamma\gamma^*}(Q^2)$. With this finding in mind we explore the idea of IVMD and the Regge models. Regge models with (infinitely many) tree-level meson and glueball exchanges are a realization of the large-N_c limit which, unlike the current quark models, incorporate confinement and comply with quark-hadron duality by generating meromorphic two-point functions for color singlet currents.

The simplest IVMD model which comes to mind has just one state saturating the subtracted dispersion relation for $F_{\pi^0\gamma\gamma^*}$. Then

$$F_{\pi^0\gamma\gamma^*}(Q^2) = \frac{1}{4\pi^2 f_\pi}\left[1 - c\frac{Q^2}{M_V^2+Q^2}\right]. \tag{6.1}$$

The fit to the combined CELLO [55], CLEO [56], and BaBar [1] data yields $c = 0.986(2)$, $M_V = 748(14)$ MeV, with $\chi^2/\mathrm{DOF} = 0.7$, thus $(1-c)$ is significantly different from 0 although numerically small. On the other hand, the fit to the CLEO data only gives $c = 0.998(18)$, $M_V = 777(44)$ MeV, $\chi^2/\mathrm{DOF} = 0.54$, with $(1-c)$ compatible with 0, or complete VMD. The result of the fit to the combined data is shown in the left panel of Fig. 3, displaying a remarkable agreement in the whole experimentally available range. The radius squared

$$b_\pi = -\left[\frac{1}{F_{\pi^0\gamma\gamma^*}(Q)}\frac{d}{dQ^2}F_{\pi^0\gamma\gamma^*}(Q)\right]\Big|_{Q^2=0} \tag{6.2}$$

becomes $b_\pi = \frac{c}{M_V^2} = 1.76(7)\ \mathrm{GeV}^{-2}$, compared to the PDG value $b_\pi = (1.76\pm0.22)\mathrm{GeV}^{-2}$. The constrain from the rare Z^0 decay is satisfied comfortably, as $|F_{Z\to\pi^0\gamma}(M_Z^2)/F_{Z\to\pi^0\gamma}(0)| = 0.014(2)$, an order of magnitude less than 0.17 of Eq. (2.7).

Next, we consider radial Regge models, with $M_n^2 = M_V^2 + an$, recalling that the large-N_c QCD involves tree-level diagrams with infinitely many states, including the radial excitations. A particular realization, the Veneziano-Dominguez model [57, 58], allows to control the asymptotic behavior of the form factor with a single parameter, b:

$$\begin{aligned} F_{\pi^0\gamma\gamma^*}(t) &= \frac{1}{4\pi^2 f_\pi} f_b(t), \\ f_b(t) &= \frac{1}{B(b-1,\frac{M_V^2}{a})}\sum_{n=0}^{\infty}\frac{\Gamma(2-b+n)}{\Gamma(n+1)\Gamma(2-b)}\frac{1}{M_n^2-t}, \end{aligned} \tag{6.3}$$

where $f_b(0) = 1$, $f_b(t=-Q^2) \sim (Q^2)^{1-b}$. The same approach was used in [59] to successfully describe the pion charge form factor, also at large Q^2. Note that $b \le 1.5$ complies to TW II, while $b > 1.5$ violates the bound. In the right panel of Fig. 3 we show the results of several Regge models: the Veneziano-Dominguez model with the fitted value $b = 1.81$ (dashed line), the Veneziano-Dominguez model with the first pole separated and fixed $b = 1.5$ (dotted line), and the subtracted Regge model (solid line). The details can be found in [32]. We note that the models differ only at very large values of Q^2, where the present experimental uncertainties are large.

Hence, it is possible to explain the BaBar data within the Regge approach, both in the TW II violating and TW II conserving versions.

Along the lines of the discussion around Eq. (3.1), another adventurous scenario can be proposed, which lies between the one mass state case and the infinitely many uniformly spaced squared mass states. We can then interpret the constant c for the IVMD ansatz of Eq. (6.1) as corresponding to a high-mass state, M_H, such that for $M_V^2 \ll Q^2 \ll M_H^2$ one has a Q^2-independent behavior. This can be represented by the simple functional form

$$F_{\pi^0\gamma\gamma^*}(-Q^2) = \frac{1}{4\pi^2 f_\pi}\left[c\frac{M_V^2}{M_V^2+Q^2} + (1-c)\frac{M_H^2}{M_H^2+Q^2}\right], \tag{6.4}$$

which resembles the IVMD ansatz of Eq. (6.1) in the range $M_V^2 \ll Q^2 \ll M_H^2$. A direct fit to the joint CLEO and BaBar data, with fixed $M_V = 748$ MeV, yields $c = 0.085(2)$ and the lowest bound (within the 95% confidence level) for the high-mass state is $M_H \sim 10$ GeV.

7. Photon momentum asymmetry

The last issue we wish to discuss in this talk is the influence of the photon momentum asymmetry parameter, A. In the BaBar kinematic setup $-q_1^2 < 0.6\ \mathrm{GeV}^2$ and $-q_2^2 > 3\ \mathrm{GeV}^2$, therefore

$$|A| = \left| \frac{q_1^2 - q_2^2}{q_1^2 + q_2^2} \right| \sim 0.9 - 0.97 \neq 1. \tag{7.1}$$

This turn out to be significant for precision fits and the optimum values of their physical parameters. In particular, the single-state IMVD model becomes

$$F_{\pi^0\gamma\gamma^*}(-Q^2) = \frac{1}{4\pi^2 f_\pi} \left[1 - c \left(1 - \frac{4M_V^4}{4M_V^4 + 4M_V^2 Q^2 + (1-A^2)Q^4} \right) \right]. \tag{7.2}$$

The fits to the combined data yield for $A = 1$, 0.975, and 0.95, respectively, $c = 0.986$, 0.978, 0.974, $M_V = 748$, 754, 768 MeV. In particular, the value of M_V increases toward m_ρ as A decreases.

8. Conclusions

These are our main conclusions:

- The second Terazawa-West bound, following from assumptions for the real part of the pion transition form factor, need not be fulfilled in field-theoretic approaches. An explicit counter-example with the Georgi-Manohar model with $g_A \neq 1$, where the chiral anomaly is fulfilled but $F_{\pi^0\gamma\gamma^*}$ tends to a constant at $t \to -\infty$. If this feature holds in QCD, it opens a possibility of solving the BaBar puzzle within models incorporating the incomplete vector-meson dominance.

- The coefficient in the TW II is large, extending the bound an order of magnitude above the data. Thus, even if it holds in the absence of polynomial contributions to the real part, it is completely ineffective for the presently-available momentum range.

- An additional constraint on the models in the time-like region of momenta follows from the rare $Z \to \pi^0\gamma$ decay. We use this bound in our considerations. For the considered models it is comfortably satisfied.

- The simplest model realizing the incomplete vector-meson dominance, including a single vector-meson state, is capable of reproducing the data for $F_{\pi^0\gamma\gamma^*}$ in the whole available experimental range, $0 < Q^2 < 35\ \mathrm{GeV}^2$.

- Within the Regge approach with infinitely many radially excited states, the data can be fitted both satisfying or violating the TW II bound.

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- The precise numerical fits are sensitive to the photon momentum asymmetry parameter, A, which affect the optimum values of the physical parameters, such as the vector meson mass. Since A is not strictly 1 in the experimental setup, the effects of kinematic cuts should be considered in precision analyses.

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