QCD-oriented nondiagonal GVDM

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Abstract. The nondiagonal generalized vector dominance model (GVDM) of photoabsorption is elaborated using QCD-motivated picture of the \( \gamma - q\bar{q} \) transition and subsequent meson dominated scattering of the \( q\bar{q} \)-pair on the nucleon. The relativistic constituent quark model for a description of the meson \( q\bar{q} \)-wave functions is used. The meson-nucleon scattering is calculated in the two-gluon exchange approximation. It is shown that the destructive interference effects and corresponding cancellations in the photoabsorption cross section formula are small, so the GVDM predictions are incorrect if no extra cut-off factors in GVDM formulas are introduced.

Key words: vector dominance, photoabsorption, hadronic scattering, constituent quarks.

1. Introduction.

According to GVDM [1] the imaginary part of the transverse forward Compton scattering amplitude (or the transverse photon absorption cross section) can be expressed in a form of the mass dispersion relation,

\[
\sigma_T(Q^2, s) = \int \frac{\rho_T(m^2, m'^2, s)m^2m'^2}{(m^2 + Q^2)(m'^2 + Q^2)} dm^2 dm'^2. \tag{1}
\]

The spectral weight function \( \rho_T \) is given by the formula of GVDM (in zero-width approximation):

\[
\rho_T(m^2, m'^2, s) = \sum_{n,n'} \delta(m^2 - m_n^2)\delta(m'^2 - m_{n'}^2) \frac{f_n}{f_{n'}} \frac{e^T_{nn'}(s)}{s}. \tag{2}
\]

Here, \( m_n \) is the vector meson mass, \( f_n \) is the meson-photon coupling constant, \( T_{nn'}(s) \) is an amplitude for the forward meson-nucleon scattering,

\[
V_n + N \longrightarrow V_{n'} + N. \tag{3}
\]

The main problem, of course, is the description of \( \rho_T(m^2, m'^2, s) \) in a region of large vector meson masses (\( m_n >> m_\rho, m_\omega, m_\phi \)). A correct treatment of
the heavy masses would provide, in particular, the convergence of the integral (1). The information about vector meson properties in the heavy mass region is rather scarce, therefore, in pre-QCD era, for the proper choice of the mass dependence of $\rho_T$ the motivation based on parton models was used. It was shown [1] that, in general, $\rho_T$ can be chosen in the form compatible with Bjorken scaling. In particular, in diagonal approximation, when

$$\rho_T(m^2, m'^2, s) = \delta(m^2 - m'^2)\rho_T(m^2, s),$$

one needs, for this compatibility, the hadronic state continuum term in $\rho_T(m^2, s)$. If the $\rho_T$-function is nondiagonal, the scaling behavior of $Q^2\sigma_T$ is possible in the more realistic case of isolated hadronic states as well. If, e.g., the $\rho_T$-function contains large negative off-diagonal contributions, scaling can be achieved through the destructive interference effects, i.e. through the strong cancellations of diagonal and off-diagonal contributions in the integral (1) (such a picture was confirmed by the direct calculation [2] of $\rho_T(m^2, m'^2, s)$ in a framework of the covariant parton model). The nondiagonal GVDM [3-5] based on this picture was really very successful in a description of the nucleon structure functions at small $Q^2$, besides, the qualitatively correct $Q^2$-dependence of the nuclear shadowing was obtained [6]. It was shown recently [7] that this model even predicts, similarly to the parton model, the color transparency effects. One should note, however, that the choice of the nondiagonal elements of $T_{nn'}$-matrix in this model has in fact no connection with the predictions of hadronic models. In this sense nondiagonal GVDM uses some fictitious vector mesons. But the original GVDM's idea is that a photon transforms virtually just into the genuine hadron states (those observed in $e^+e^-$-annihilation) which subsequently scatter from the target nucleon. Correspondingly, the $f_n$-constants in Eq.(2) are expressed through the leptonic widths of these states. The logic of GVDM should be such that the hadronic physics is a starting point and the scaling behavior of $Q^2\sigma_T$ is the (nonnecessary, in principle) consequence.

In last few years many works appeared applying the diagonal approximation for $\rho_T$. Some of them introduce new parametrizations for $\rho_T(m^2, s)$-dependencies based solely on the experimental data, others sharply cut off the heavy mass tail of the meson mass spectrum adding, instead, the large softly interacting $\bar{q}q$-component [8] or the large direct component [9].

The common feature of all these recent works is, again, a lack of the attention paid to the purely hadronic aspects of the problem. Two (at least) main questions should be studied: i) how important are nondiagonal elements of $T_{nn'}$ in the integral (1) and ii) how the problem of the heavy meson masses is resolved (in particular, is there a necessity in some extra cut-off factor in GVDM formulas).

In the present paper we try to answer both these questions using relativistic constituent quark model of vector mesons and the two-gluon exchange approximation for a calculation of the meson-nucleon scattering amplitudes.
A correct use of the hadronic basis is very important for a quantitative
description of the nuclear shadowing (which is quite sensitive just to a space-
time picture of the process) and, especially, for the precise calculation of $\sigma_{\gamma N} = \sigma_T(s, 0)$ at very high photon energies (at the energy region which is of much interest for astrophysics and cosmic ray physics and which is experimentally inaccessible).

Concluding, one should add that now there are approaches [10,11] in which
the hadronic basis in the photoabsorption description at small $Q^2$ is completely
discarded. The price for this is a necessity of an use of new parameters which
either have a rather dubious physical sense (as the virtuality depe-
ning constituent quark mass [10,11]) or are badly known (as the confinement radius
[11]).

2. The model of the hadronic amplitudes.

We will use the simplest model of VN-scattering: two-gluon exchange
approximation. For a calculation of the corresponding diagrams one must know, in
particular, $q\bar{q}$-wave functions of the mesons. In a relativistic constituent quark
model these wave functions are obtained from the Bethe-Salpeter(BS) equation
(we consider the case of scalar identical quarks):

$$i(2\pi)^4 \Phi(P, q) = \frac{1}{\Delta_1 \Delta_2} \int d^4 q' K(q - q') \Phi(P, q').$$

Here, $\Phi(P, q)$ is the BS wave function, $P = k_1 + k_2, q = \frac{1}{2}(k_1 - k_2), k_{1,2}$ are
quark 4-momenta, $\Delta_{1,2} = k_{1,2}^2 - m_q^2; K(q - q')$ is the $q\bar{q}$-interaction kernel.

For the utilization of this equation it is convenient to use the quasipotent-
tial formalism in a light-front form (see, e.g., [12,13]). Variables needed for a
description of an internal motion of the constituents in this form of dynamics
are light-front momenta $\rightarrow k_1, \rightarrow k_2 (k = \rightarrow k_1, \rightarrow k_2)$ which transform covariantly under
the kinematic Lorentz transformations. Correspondingly, $q_{\perp}$-component of the
relative momentum $q$ is restricted. The choice of a concrete form of this re-
striction is not unique, however. The most simple and natural way is, in our
opinion, an use of the covariant condition $Pq = 0$ [13]. In this case one has (if
$P_+, P_\perp = M^2, P_\perp = 0; M$ is the bound state mass)

$$q_+ = -\frac{P_+ q_+}{P_+} = -\frac{M^2 q_+}{P_+^2}; q^2 = -q_\perp^2 - \frac{M^2 q_+^2}{P_+^2} = -q_\perp^2 - M^2 y^2.$$  

(6)

Now we can introduce 3-dimensional(3D) "inner momentum" $\tilde{q}$ ($\tilde{q}_\perp = q_\perp; \tilde{q}_3 = M y$) which is, according to Eqs.(6), an argument of the interaction kernel, $K = K(q - q')$. The corresponding reduction of the BS-equation can be performed
by the integration both sides of Eq.(5) over \( q_- \) and using the basic formula

\[
\frac{1}{2} \int \frac{1}{\Delta_1 \Delta_2} dq_- = \frac{2\pi i}{2F_+ M^2(q_-) - q_-^2 - m_q^2}.
\]

(7)

The integration in Eq.(7) is done by contour methods and corresponds to putting one of the quarks on its mass shell. The resulting 3D equation is

\[
(\hat{q}^2 + m_q^2 - \frac{M^2}{4})\Phi(\hat{q}) = \frac{1}{16\pi^3M} \int d\hat{q} K(\hat{q} - \hat{q}')\Phi(\hat{q}').
\]

(8)

The wave function \( \Phi(\hat{q})(= \Phi(\hat{q}_\perp, y)) \) is simply connected with the Vqq-vertex function:

\[
\Gamma(\hat{q}) = \Phi(\hat{q})(\hat{q}_\perp^2 + m_q^2 - \frac{M^2}{4}).
\]

(9)

Now we are able to calculate the two-gluon exchange diagrams starting from the general 4D expressions and reducing them, with an aid of Eqs.(7,9), to the 3D form. We keep only the terms of a leading order in \( s \) and neglect the longitudinal momentum transfer. The resulting amplitude for the meson-meson scattering is

\[
F(s, t = -\frac{Q^2}{4}) = \frac{i}{(2\pi)^3} \frac{8}{9} g^4 \int \frac{d^2k_\perp}{(\frac{Q}{2} - k_\perp)^2 (\frac{Q}{2} + k_\perp)^2} \times
\]

\[
|F_1(Q^-) - F_1(4k_\perp^-)| |F_2(Q^-) - F_2(4k_\perp^-)|.
\]

(10)

Here, \( F \) is the meson formfactor given by the formula

\[
F(Q^2) = \int d^2q_\perp \int dy (1 - 4y^2) \Phi(q_\perp^-, y) \Phi(q_\perp^+ + Q^-/2, y).
\]

(11)

Similar (but not identical) expression for the amplitude was obtained long ago [14] employing the eikonal formalism.

The amplitude of the meson-nucleon scattering is obtained from Eq.(10) by the replacement:

\[
[F_2(Q^2) - F_2(4k_\perp^-)] \to \frac{3}{2} V(k_\perp, Q^-).
\]

(12)

The V-factor describes the ggNN-vertex. We estimate this factor using, for simplicity, the approach of ref.[14]:

\[
V(k_\perp, Q^-) \approx \frac{1}{3} \sum_i (e^{iQx_i}) - \frac{1}{6} \sum_{i \neq j} (e^{i[(Q + k_\perp)x_i] + (Q - k_\perp)x_j}) \approx
\]

\[
= e^{-\frac{Q^2}{2\Lambda^2}} - e^{-\frac{k_\perp^2}{2\Lambda^2}} (\frac{Q^2}{4} + 3k_\perp^2).
\]

(13)
Here, $< ... >$ denotes the expectation value in the nucleon bound state and it is assumed that the three quarks have the same Gaussian distribution; $< r^2_N >$ is the mean squared radius of the nucleon.

For the solution of Eq.(8) we assume that kernel K has only the long range confining term of the hadronic oscillation type [13]:

$$K(q - q') = (2\pi)^3 \omega^2_{qq}(\nabla_q^2 + \omega_0^{-2})\delta^3(q - q'),$$

with two parameters: $\omega^2_{qq}$ (a "spring constant") and a zero-point energy $\omega_0$.

Equation (8) with this kernel formally coincides with the equation for a quantum-mechanical 3D-oscillator:

$$\frac{\tilde{q}^2}{\omega^2_{qq}} + m^2_q - \frac{M^2}{4} \Phi(q) = \frac{\omega^2_{qq}}{2M}(\nabla_q^2 + \omega_0^{-2})\Phi(q).$$

Solutions of Eq.(15), its eigenfunctions and eigenvalues are well known. We will use them for a description of the $\rho$-family. The mass spectrum of radial excitations is:

$$1/2\beta^2(M^2/4 - m^2_q + \beta^4\omega_0^{-2}) = N + 3/2; \quad N = 0, 2, 4, ... .$$

Deriving Eq.(16) we assume that $\beta^2 \equiv \omega_q/\sqrt{2M}$ is a constant (i.e., is independent on M). In this case the meson mass spectrum has the form $m^2_n = a + bn$. With the numerical values

$m_\rho = 0.77, \quad m_{\rho'} = 1.45, \quad m_q = 0.3$

one has

$\beta^2 = 0.094 \text{ Gev}^2, \quad \omega_0^2 = 0.04 \text{ Gev}^2$

and, finally,

$m^2_n \cong m^2_\rho(1 + 2.55n); \quad n = 0, 1, 2, ... .$

The $\rho$-meson wave function is

$$\Phi_0(q) = N_0\exp[-(q^2_\perp + m^2_\rho y^2)/2\beta^2],$$

and $N_0$ is determined from Eq.(11), using the condition $F(0)=1$. In the $(r_\perp, y)$-representation one has

$$\Phi_0(r_\perp, y) = N_0\exp[-r^2_\perp/2\beta^2]exp[-m^2_\rho y^2/2\beta^2].$$

Using the $(r_\perp, y)$-space, the $V_{\rho N}$-scattering amplitude can be written as

$$F_{nn}(s, t) = \int d^2r_\perp dy F_{\rho n}(s, t)\Phi^2_n(r_\perp, y) \equiv < n \mid F_{r_\perp}^\rho(s, t) \mid n >,$$
Here, \( F_{r_\perp} \) is an "eigenamplitude", i.e. an amplitude for the scattering of the \( q\bar{q} \)-pair with a fixed \( r_\perp \) on the nucleon. Eq.(20a) could be written without any derivation, using only the simple physical fact that, due to the large lifetime of \( q\bar{q} \)-fluctuations at large \( s \) the values \( r_\perp, y \) are "frozen" in the scattering process (and therefore they are "eigenvalues" of the scattering matrix). The concrete expression for \( F_{r_\perp} \) (Eq.(20b)) is given by the model.

Integrating over asimutal angles in Eqs.(20) we reduce the problem to a calculation of \( F_{r_\perp}^{|r_\perp|} \equiv F_{r_\perp} \). This amplitude depends only on two parameters, \( \alpha_s \) and \( \mu_g \), effective gluon mass (omitted in the above-cited expressions for brevity’s sake). Going into impact parameter space, we introduce the opaque function

\[
\Omega_{r_\perp}(s, b) = \frac{1}{2\pi} \int \frac{1}{4\pi s} F_{r_\perp}(s, t) e^{iQ_b} d^2Q. \tag{21}
\]

The numerical calculation shows that \( \Omega_{r_\perp} \) can be parametrized with a large accuracy by the Regge-type expression:

\[
\Omega_{r_\perp}(s, b) = \frac{\sigma(r_\perp)}{4\pi B_{r_\perp}} e^{\frac{b^2}{4B_{r_\perp}}}, \tag{22}
\]

where

\[
\sigma(r_\perp) = \frac{1}{s} \text{Im} F_{r_\perp}(s, 0); \quad B_{r_\perp} = \frac{\sigma(r_\perp)}{4\pi \Omega_{r_\perp}(s, 0)}. \tag{23}
\]

We took in this analysis \( \mu_g = \mu_\pi \) and normalized \( \sigma(r_\perp) \) on the pion data at medium energies (\( \sqrt{s} = 10 \text{ Gev} \)), in accordance with the additive quark model relation

\[
\sigma_{pp} = \frac{1}{2}(\sigma_{\pi^+} + \sigma_{\pi^-}). \tag{24}
\]

For this normalization we used the unitarized scattering amplitude

\[
T_{pp}(s, 0) = \int < \rho | 1 - e^{-\Omega_{r_\perp}(s, b) \rho} | \rho > d^2\rho. \tag{25}
\]

Up to now in our model \( F_{r_\perp} \sim s \) so that \( \Omega_{r_\perp} \) does not depend on the energy. To take into account this dependence we modify Eq.(22) adding a new Regge-type term:

\[
\Omega_{r_\perp}(s, b) = \frac{\sigma(r_\perp)}{4\pi} \left\{ \frac{1}{B_{r_\perp}} e^{-\frac{b^2}{4B_{r_\perp}}} + \frac{1}{B_{r_\perp}} \frac{1}{R_{r_\perp}} e^{\Delta_F} \xi e^{-\frac{b^2}{2(B_{r_\perp} + 2\alpha_F')}} \right\}, \tag{26}
\]

\[
\xi = ln \frac{s}{s_0} - \frac{i\pi}{2}; \quad \Delta_F = \alpha_F - 1 > 0; \quad \alpha_F' \neq 0.
\]
Writing Eq.(26) we suppose a two-pole Regge-parametrization [15] of the opaque function; by assumption, both trajectories give at small $\xi$ the same diffraction slopes and $R$ does not depend on $r_\perp$. Three new parameters ($R, \Delta_F, \alpha'_F$) are determined from data on a $s$-dependence of hadronic total amplitudes.

Now we have all necessary for a calculation of the amplitudes. It is evident that the nondiagonal amplitudes are given by a simple generalization of Eq.(20a):

$$F_{nn'}(s, t) = <n | F_{r_\perp} (s, t) | n'> .$$

(27)

In the nondiagonal case one has, instead of the elastic scattering $V_n N \rightarrow V_n N$, the diffraction dissociation $V_n N \rightarrow V_n' N$.

3. Cut-off factors.

The first stage of the photoabsorption process is the $\gamma \rightarrow q \bar{q}$-transition. The differential probability of this transition is

$$dP_{qq} = C \frac{1}{\mu_\perp^2} \{x^2 + (1-x)^2 + \frac{2m_q^2}{\mu_\perp^2} x(1-x)\} dx dp_\perp^2 .$$

(28)

Here, $C$ is a known constant, $\mu_\perp$ is the quark transverse mass, $\mu_\perp = \sqrt{p_\perp^2 + m_q^2}$, $x$ is the fraction of the photon 3-momentum carried by the quark. An invariant mass of the $q \bar{q}$-pair is

$$M_{qq}^2 = \mu_\perp^2 / x(1-x) .$$

(29)

The last term in the curly brackets in Eq.(30) is of the order $\sim m_q^2 / M_{qq}^2$ and can be safely neglected (in this section $m_q$ is the current quark mass). Going over from $(p_\perp^2, x)$- to $(M_{qq}^2, x)$-variables in Eq.(28) we obtain

$$dP_{qq} \approx C \frac{1}{M_{qq}^2} [x^2 + (1-x)^2] dx dM_{qq}^2 .$$

(30)

It is easy to show that the average transverse size of the $q \bar{q}$-pair is given by the formula

$$\overline{r}_\perp = v_{\perp, relative} \cdot \tau_F \approx \frac{p_\perp}{\mu_\perp} (1 + \frac{Q^2}{M_{qq}^2})^{-1} \frac{m_q = Q^2 = 0}{1} \frac{1}{p_\perp} .$$

(31)

In the present paper we consider only the case $Q^2 = 0$. Evidently, $\overline{r}_{\perp, max} \sim \sqrt{m_q^{-1}}$ is much larger than a typical transverse size of hadrons. Our basic assumption is the following: an interaction of the $q \bar{q}$-pair with the nucleon is meson-dominated if (and only if) this pair is wide enough (i.e. if $\overline{r}_\perp > \overline{r}_h^0 \sim \sqrt{<r_h^2>}$; only in
this case confinement forces are effective and pull the pair’s particles together. It is known that the narrow pairs weakly interact with a nucleon (due to the color transparency phenomenon). Therefore, the corresponding nonVDM contribution is, fortunately, small and can be taken into account by a slight change of the $r^0_\perp$-parameter. Very narrow pairs (with $p_{\perp} \geq 2$ Gev) interact with a nucleon purely pertubatively and must be considered separately (the corresponding “anomalous” contribution to $\sigma_{\gamma N}$ is essential only at very high energies).

The restriction of the $q\bar{q}$-pair’s phase volume was discussed by many authors beginning from its suggestion in ref.[16]. It follows from Eqs.(29,30) that, at fixed $M^2_{q\bar{q}}$, the relative part of pair’s phase volume having $p_{\perp}$ in the limits ($\sim m_q / p_{\perp}^{\text{max}}$) is given by

$$\eta \approx 3 \left( \frac{p_{\perp}^{\text{max}}}{M_{q\bar{q}}} \right)^2, \text{ for } M^2_{q\bar{q}} \gg (p_{\perp}^{\text{max}})^2. \quad (32)$$

Using Eq.(31) we introduce the restricting factors. Identifying $M_{q\bar{q}}$ with $m_n$, we have for each vector meson:

$$p_{\perp n}^{\text{max}} = (\alpha \sqrt{< r^2_n >})^{-1}; \quad \eta_n \approx 3 \left( \frac{p_{\perp n}^{\text{max}}}{m_n} \right)^2. \quad (33)$$

We assume that parameter $\alpha$ is the same for all mesons. Evidently, the restriction is absent if $M_{q\bar{q}} < 2p_{\perp}^{\text{max}}$ (it appears that this inequality is valid for $\rho, \omega$-mesons only).

Finally, in GVDM formulas the following substitutions must be done:

$$e \rightarrow e \sqrt{\eta_n} \equiv e \frac{f_n}{f_n} \quad (34)$$

The mean square transverse radii of vector mesons can be calculated using the wave functions described in the previous section.

4. Results and conclusion.

Our formula for the photoabsorption cross section is ($Q^2 = 0$)

$$\sigma_{\gamma N} = \sum_{n,n'} \frac{e^2}{f_n f_n'} \frac{\text{Im} T_{nn'}}{s} \equiv \sum_{n,n'} \sigma_{\gamma N}^n. \quad (35)$$

The coupling constants $f_n$ are simply connected with the lepton widths $\Gamma_n (V_n \rightarrow e^+ e^-)$. In principle, these widths should be calculated with an aid of the quasipotential formalism used above. It will be done in a separate paper. Now we assume, as usual, that $\Gamma_n \sim m_n^{-1}$ (here $n$ is, as earlier, the meson number in the family). From this it follows that $f_n \sim m_n$. 

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In the table 1 we present our results for the $\rho$-family contribution to $\sigma_{\gamma N}$ ($\sqrt{s} = 10\text{Gev}$). Each number in the second line of the table is one of the terms in the sum of Eq.(35) calculated without cut-off factors. The basic coupling constant is well known, $f_\rho^2/4\pi = 2.25$. The third line of the table contains $\sigma^{nn'}_{\gamma N}$ obtained after an insertion of the cut-off factors, $f_n \rightarrow \tilde{f}_n$.

| Parameter | $\rho \rho'$ | $\rho' \rho''$ | $\rho' \rho''$ | $\rho \rho'$ | $\rho' \rho''$ | $\rho' \rho''$ |
|-----------|---------------|-----------------|-----------------|---------------|-----------------|-----------------|
| $\sigma^{nn'}_{\gamma N}$ without cut-off factors | 77.9 | 41.86 | 28.49 | 12.9 | -5.6 | -5.04 |
| $\sigma^{nn'}_{\gamma N}$ with cut-off factors | 77.9 | 10.56 | 2.66 | 6.47 | -1.71 | -0.77 |

Parameter $\alpha (= 0.474)$ was found by the comparison of the theoretical $\sigma_{\gamma N}$ with the experimental value ($\sim 115 \mu bn$). The contribution of $(\omega, \phi)$-families was estimated in diagonal approximation with the result:

$$\sum_{n(\omega, \phi)} \sigma^{nn}_{\gamma N} = 8.65 + \frac{2.3}{\alpha^2}. \quad (36)$$

The energy dependence of $\sigma_T$ is shown on fig.1. The following values of the parameters were used:

$$R = 22; \quad \Delta_F = 0.25; \quad \alpha_F = 0.13.$$

The numerical results obtained in the present model lead to the following conclusions.

1. If no cut-offs are introduced, GVDM is not able to describe photoabsorption data. Even the simplest variant of the GVDM containing only $\rho$ and $\rho'$ give too large value of $\sigma_{\gamma N}$. Nondiagonal contributions are not negligibly small. Destructive interference effects proposed in [3-5] are not effective (in particular, the largest nondiagonal term ($\rho\rho'$) is positive).

2. The introduction of the cut-off factors motivated by QCD can give the correct predictions. This is reached without nonnatural break of the meson mass spectrum (the value of the heaviest mass in GVDM expressions is determined solely by the condition that the longitudinal size of the fluctuation must exceed the target size). In the present model only one parameter ($\alpha$) is needed for the description of all cut-offs (despite the fact that $p_{max}^\parallel$ values are different for different mesons). In the scheme with the cut-offs the nondiagonal contributions have the same order of magnitude as neighbouring diagonal ones, so GVDM developed in the present paper is an essentially nondiagonal model.
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