factorization will not go through. The only Regge-pole property left is the asymptotic behavior in ν.

[5]A. M. Boyarski, F. Bulos, W. Busza, R. Diebold, S. D. Ecklund, G. E. Fischer, J. R. Rees, and B. Richter, Phys. Rev. Letters 20, 300 (1968).

[6]N. F. Bali, G. F. Chew, and A. Pignotti, Phys. Rev. Letters 19, 614 (1967). See also J. Finkelstein and K. Kajantie, Phys. Letters 26B, 305 (1968).

RADIATIVE MESON DECAYS IN BROKEN SU(3)*

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We consider vector and pseudoscalar meson decays using vector gauge fields with current mixing. The processes $\omega\to 3\pi$, $\omega\to\pi\gamma$, and $\pi\to 2\gamma$ are fitted with one value of the gauge field's coupling g, which also determines leptonic decay rates of vector mesons. Octet breaking of the underlying VVP strong interaction introduces four parameters. The available experimental information on π, η, ω, and φ decays gives relations among these parameters which predict rates for decays such as $K^*\to K\gamma$.

The vector-dominance approach to meson decays[1] gives a good qualitative account of a number of strong and radiative decays, but the unbroken SU(3) version of this approach has certain quantitative failures. Examples are the ratios $\Gamma(\omega\to\pi\gamma)/\Gamma(\omega\to 3\pi)$ and $\Gamma(\eta\to\pi\pi\gamma)/\Gamma(\eta\to 2\gamma)$ which differ markedly from the predicted values.[2] Another experimental disagreement with the unbroken SU(3) prediction, which has been recently established,[3] is the ratio $\Gamma(\eta\to 2\gamma)/\Gamma(\pi^0\to 2\gamma)$, the disagreement in this ratio being about six. We study these and other decays here, using octet breaking as well as some recent improvements in the theory of vector gauge fields within the effective-Lagrangian framework.[4]

The fundamental strong-interaction term responsible for the processes we study has the form

$$\mathfrak{L}_{PVV}=\tfrac{1}{4}\epsilon_{\alpha\beta\mu\nu}(hD^{abc}V_{\alpha\beta}{}^{a}V_{\mu\nu}{}^{b}P^{c}+\lambda D^{ab}V_{\alpha\beta}{}^{a}P^{b}V_{\mu\nu}{}^{0}),\qquad(1)$$

where $V_\mu{}^0$ is an SU(3)-singlet vector meson, and $V_\mu{}^a$ and P^b are, respectively, vector and pseudoscalar octets: $a,b,c=1,\cdots,8$. In octet-broken SU(3) the D's have the general form[5]

$$D^{abc}=d^{abc}+\sqrt{3}\epsilon_1 d^{abd}d^{d8c}+\tfrac{1}{2}\sqrt{3}\epsilon_2(d^{acd}d^{d8b}+d^{bcd}d^{d8a})+(\epsilon_3/\sqrt{3})\delta^{ab}\delta^{c8},\qquad(2)$$

$$D^{ab}=\delta^{ab}+\sqrt{3}\epsilon_4 d^{ab8}.\qquad(3)$$

The vector fields are described by the generalized Yang-Mills Lagrangian

$$\mathfrak{L}_V=-\tfrac{1}{4}K^{ab}V_{\mu\nu}{}^{a}V_{\mu\nu}{}^{b}+\tfrac{1}{2}m^2V_\mu{}^aV_\mu{}^a-\tfrac{1}{4}K^{00}V_{\mu\nu}{}^{0}V_{\mu\nu}{}^{0}+\tfrac{1}{2}m^2V_\mu{}^0V_\mu{}^0-\tfrac{1}{2}K^{80}V_{\mu\nu}{}^{8}V_{\mu\nu}{}^{0},\qquad(4)$$

where

$$V_{\mu\nu}{}^{a}=\partial_\mu V_\nu{}^a-\partial_\nu V_\mu{}^a-gf^{abc}V_\mu{}^bV_\nu{}^c$$

and where the K's are responsible for the observed mass splittings among the nine vector mesons and for the ω-φ mixing in the vector-mixing model[6-8]; K^{ab} is the diagonal matrix

$$K^{ab}=\delta^{ab}+\sqrt{3}\epsilon_0 d^{ab8}.\qquad(5)$$

Diagonalizing (4) in terms of physical ω and φ, we obtain

$$V_\mu{}^{1,2,3}=\frac{1}{\sqrt{K_\rho}}\rho_\mu{}^{1,2,3};$$

$$V_\mu{}^{4,5,6,7}=\frac{1}{\sqrt{K_{K^*}}}K_\mu{}^{*4,5,6,7};$$

$$V_\mu{}^{8}=-\frac{\sin\theta}{\sqrt{K_\omega}}\omega_\mu+\frac{\cos\theta}{\sqrt{K_\varphi}}\varphi_\mu;$$

$$V_\mu{}^{0}=\frac{\cos\theta}{\sqrt{K_\omega}}\omega_\mu+\frac{\sin\varphi}{\sqrt{K_\varphi}}\varphi_\mu,\qquad(6)$$

with

$$K_i=m^2/m_i{}^2\quad(i=\rho,K^*,\omega,\varphi),$$

and[9]

$$m = 847\text{ MeV},\quad \theta = 27.5°.$$

Following the formalism of Refs. 7 and 8, we get the effective electromagnetic interaction

$$\mathfrak{L}_{\text{em}} = (em^2/g)[K_\rho^{-1/2}\rho_\mu^{\ 3} - (3K_\omega)^{-1/2}\omega_\mu \sin\theta + (3K_\varphi)^{-1/2}\varphi_\mu \cos\theta]A_\mu. \quad (7)$$

Using expressions (1), (6), and (7) we obtain amplitudes for the processes listed in Table I. These differ from amplitudes previously given in the literature by symmetry-breaking factors displayed in curly brackets which approach unity in the SU(6)-symmetric limit.[10] For example, for $\eta \to \pi^+\pi^-\gamma$ we calculate the invariant amplitude

$$F\epsilon_{\alpha\beta\mu\nu} e_\alpha^{(\gamma)} k_\beta^{(\gamma)} p_\mu^{(1)} p_\nu^{(2)} \times (p^2 - m_\rho^2)^{-1} g_{\rho\pi\pi}(p^2)/g_{\rho\pi\pi}(0), \quad (8)$$

with $p_\mu^{(i)}$ the pion momenta and $p_\mu = p_\mu^{(1)} + p_\mu^{(2)}$. The three diagrams contributing to the $\omega \to 3\pi$ amplitude have the same form as (8) with superscript (γ) replaced by (ω). The factor $g_{\rho\pi\pi}(p^2)/g_{\rho\pi\pi}(0)$ is the $\rho\pi\pi$ form factor normalized at $p^2 = 0$, with

$$g_{\rho\pi\pi}(0) = g\sqrt{K_\rho}. \quad (9)$$

Similarly, the invariant amplitudes for the two-body decays listed in Table I have the form

$$F\epsilon_{\alpha\beta\mu\nu} e_\alpha^{(V_1)} k_\beta^{(V_1)} e_\mu^{(V_2)} k_\nu^{(V_2)}, \quad (10)$$

where the labels (V_i) stand for γ or vector meson.

In arriving at the expressions in Table I we fixed the value of λ so that $\mathfrak{L}_{PVV}$ gives no contribution to the $\varphi \to 3\pi$ decay amplitude, which is known to be very small. This gives

$$\lambda = -\frac{2h\cot\theta}{\sqrt{3}}\left(\frac{1+\epsilon_1}{1+\epsilon_4}\right). \quad (11)$$

Consider now the processes $\omega \to 3\pi$, $\omega \to \pi\gamma$, and $\pi^0 \to 2\gamma$ which depend only on $h(1+\epsilon_1)$ and on different powers of $g^2/4\pi$, as displayed in Table II. Comparison of the calculated decay widths with the observed ones[11] overdetermines g and provides a consistency test. Agreement is obtained with

$$g^2/4\pi = 3.35,\quad (m_\pi^{\ 2}h^2/4\pi)(1+\epsilon_1)^2 = 0.10. \quad (12)$$

This is remarkable since previous authors, using unbroken SU(3), have found discrepancies of factors of 2 or more in comparing these decays with each other. The improvement in our treatment is due solely to the use of vector mixing. {Our determination of $g^2/4\pi = 3.35$ [Eq. (12)] permits an evaluation of absolute rates for the leptonic decays of vector mesons. We find[12] $\Gamma(\rho \to l^+l^-) = 5.0$ keV, $\Gamma(\varphi \to l^+l^-) = 1.0$ keV, $\Gamma(\omega \to l^+l^-) = 0.36$ keV. The values for the absolute rates given in Ref. 7 for ω and φ decays into leptons are twice as large as ours. This is because those authors use a value of $g^2/4\pi$ derived from the strong decay $\varphi \to K\bar{K}$.}

In calculating $\omega \to 3\pi$ we set $g_{\rho\pi\pi}(p^2) = g_{\rho\pi\pi}(0)$. With the same assumption, namely, neglect of the momentum dependence of the form factor, the coupling constant of the $\rho \to \pi\pi$ decay would be from (9)

$$g_{\rho\pi\pi}^{\ 2}(0)/4\pi = 2.76 \pm 0.55, \quad (13)$$

corresponding to a rho width of 144 ± 30 MeV, in reasonable agreement with experiment. We also predict

$$\Gamma(\rho \to \pi\gamma) = 0.075\text{ MeV}, \quad (14)$$

about 60% of the unbroken-symmetry value.

We emphasize that our results so far are independent of the dynamical breaking exhibited in Eqs. (2) and (3), that is, independent of the choice of the ϵ_i. Thus, the quality of the agreement lends support to the validity of the vector-mixing vector-dominance model.[13] In contrast to this, the observed ratio[3,11] $\Gamma(\eta \to 2\gamma)/\Gamma(\pi^0 \to 2\gamma)$ is larger than the SU(3)-symmetric prediction by a factor of 6, and this situation is unaltered by vector mixing. A similar discrepancy holds for the branching ratio $\Gamma(\eta \to \pi^+\pi^-\gamma)/\Gamma(\eta \to 2\gamma)$. Using our model with $g_{\rho\pi\pi}(0)$ given by Eq. (13) and with all ϵ_i set equal to zero, we obtain 0.50 for this ratio compared with the experimental ratio[11] of 0.14, so that vector mixing alone is not sufficient.[14]

Using the matrix elements (8) and (10) with F given in Table I, a comparison with the experimental rates gives

$$(1-\epsilon_1+2\epsilon')^2 = 6.0(1+\epsilon_1)^2 \quad (15)$$

and

$$(1-\epsilon_1+\epsilon')^2 = 1.7(1+\epsilon_1)^2, \quad (16)$$

with $\epsilon' = \epsilon_2 + \epsilon_3$. The numbers 6.0 and 1.7 are determined by the experimental rates to about 20%

Table I. Broken-SU(3) amplitudes in the vector-meson-dominance model.

Process	F
$\omega \to 3\pi$	$4hg\left\{(1+\epsilon_1)/K_\rho\sqrt{3K_\omega}\sin\theta\right\}$
$\varphi \to 3\pi$	0 (input)
$\varphi \to \pi\gamma$	0
$\omega \to \pi\gamma$	$-\frac{2he}{g}\left\{\frac{1+\epsilon_1}{\sqrt{3K_\omega}\sin\theta}\right\}$
$\rho \to \pi\gamma$	$\frac{2he}{3g}\left\{\frac{1+\epsilon_1}{\sqrt{K_\rho}}\right\}$
$\varphi \to \eta\gamma$	$-\frac{4\sqrt{2}\,he}{3\sqrt{3}\,g}\left\{\frac{\sqrt{3/2}\cos\theta}{2\sqrt{K_\varphi}}\left[1-\epsilon_1-\epsilon_2-\epsilon_3+\frac{1+\epsilon_1}{1+\epsilon_4}(1-\epsilon_4)\right]\right\}$
$\omega \to \eta\gamma$	$-\frac{2he}{3\sqrt{3}\,g}\left\{\frac{\sqrt{3}}{\sqrt{K_\omega}\sin\theta}\left[-\sin^2\theta\,(1-\epsilon_1-\epsilon_2-\epsilon_3)+\cos^2\theta\,\frac{(1+\epsilon_1)(1-\epsilon_4)}{1+\epsilon_4}\right]\right\}$
$\rho \to \eta\gamma$	$\frac{2he}{\sqrt{3}\,g}\left\{\frac{1}{\sqrt{K_\rho}}(1-\epsilon_1+\epsilon_2+\epsilon_3)\right\}$
$K^{*+} \to K^+\gamma$	$\frac{2he}{3g}\left\{\frac{1}{\sqrt{K_{K^*}}}\left(1-\frac{\epsilon_1}{2}+\frac{3\epsilon_2}{4}\right)\right\}$
$K^{*0} \to K^0\gamma$	$-\frac{4he}{3g}\left\{\frac{1}{\sqrt{K_{K^*}}}\left(1-\frac{\epsilon_1}{2}\right)\right\}$
$\eta \to \pi^+\pi^-\gamma$	$\frac{4he}{\sqrt{3}}\left\{\frac{1}{K_\rho}(1-\epsilon_1+\epsilon_2+\epsilon_3)\right\}$
$\eta \to 2\gamma$	$\frac{4he^2}{3\sqrt{3}\,g^2}\left\{1-\epsilon_1+2\epsilon_2+2\epsilon_3\right\}$
$\pi^0 \to 2\gamma$	$\frac{4he^2}{3g^2}\left\{1+\epsilon_1\right\}$

Table II. Processes depending only on g and $h(1+\epsilon_1)$.

Process	Γ_{calc} (MeV)	Γ_{obs}[a] (MeV)
$\omega \to 3\pi$	$\left(\frac{h^2 m_\pi^{\ 2}}{4\pi}\right)\left(\frac{g^2}{4\pi}\right)(1+\epsilon_1)^2(32.2)$	11.0 ± 1.3
$\omega \to \pi\gamma$	$\left(\frac{h^2 m_\pi^{\ 2}}{4\pi}\right)\left(\frac{g^2}{4\pi}\right)^{-1}(1+\epsilon_1)^2(38.4)$	1.16 ± 0.24
$\pi^0 \to 2\gamma$	$\left(\frac{h^2 m_\pi^{\ 2}}{4\pi}\right)\left(\frac{g^2}{4\pi}\right)^{-2}(1+\epsilon_1)^2(813\times 10^{-6})$	$(7.4 \pm 1.5)10^{-6}$

[a]See Ref. 11.

accuracy. We get the following solutions:

$$\epsilon_1 = 0.77, \quad \epsilon' = 2.1; \tag{17a}$$

$$\epsilon_1 = 1.3, \quad \epsilon' = -2.7; \tag{17b}$$

$$\epsilon_1 = -1.5, \quad \epsilon' = -1.9; \tag{17c}$$

$$\epsilon_1 = -0.67, \quad \epsilon' = -1.2. \tag{17d}$$

Each solution leads to a predicted rate for $K^{*0} \to K^0\gamma$ (in MeV):

$$0.028, \tag{18a}$$

$$0.0055, \tag{18b}$$

$$2.9, \tag{18c}$$

$$3.8. \tag{18d}$$

On the other hand, the $K^{*+} \to K^+\gamma$ decay rate depends on the undetermined parameter ϵ_2, as well as on ϵ_1. If $\epsilon_2 = 0$, $\Gamma(K^{*0} \to K^0\gamma)/\Gamma(K^{*+} \to K^+\gamma) = 4$, the symmetry limit, for any choice of ϵ_1. Measurement of these two rates will select one of the four solutions for ϵ_1 and ϵ' and will also determine ϵ_2.

We can, in fact, already exclude the solutions (c) and (d) by considering the experimental upper limits for the decay modes $\varphi \to \eta\gamma$ and $\omega \to \eta\gamma$,[15] and at the same time determine the allowed range of the parameter ϵ_4. We obtain, namely,

$$\Gamma(\varphi \to \eta\gamma) = (0.098\ \mathrm{MeV})\left(\frac{1-\epsilon_1-\epsilon'}{1+\epsilon_1} + \frac{1-\epsilon_4}{1+\epsilon_4}\right)^2 < 0.27\ \mathrm{MeV} \tag{19}$$

and

$$\Gamma(\omega \to \eta\gamma) = (0.058\ \mathrm{MeV})\left(-0.214\frac{1-\epsilon_1-\epsilon'}{1+\epsilon_1} + 0.790\frac{1-\epsilon_4}{1+\epsilon_4}\right)^2 < 0.18\ \mathrm{MeV}. \tag{20}$$

For solution (a) these inequalities determine $-0.3 < \epsilon_4 < 6$; for solution (b) the allowed range is $\epsilon_4 > 0.2$ or $\epsilon_4 < -3.3$. We note also that in the case of solution (a), $|\epsilon_4| \gg 1$ gives $\Gamma(\omega \to \eta\gamma) = 0.021$ MeV and $\Gamma(\varphi \to \eta\gamma) = 0.39$ MeV, which does not greatly violate the quoted experimental upper limit. On the other hand, solutions (c) and (d) are incompatible with the pair of inequalities.

With regard to $\Gamma(\rho \to \eta\gamma)$, this rate is in our scheme entirely determined by the measured rate $\Gamma(\eta \to \pi^+\pi^-\gamma)$, as can be seen from Table I. Its predicted value is

$$\Gamma(\rho \to \eta\gamma) = 0.050\ \mathrm{MeV}. \tag{21}$$

Summarizing our results, we emphasize that once g and λ are fixed from the strong interaction processes[16] $\rho \to \pi\pi$ and $\varphi \to \rho\pi$, then the observed rates for the three processes in Table II are compatible with a single value for $h(1+\epsilon_1)$. We regard this as a significant test for the vector-mixing, vector-dominance model. We found next that the observed rates for the radiative η decays determined four possible sets of values [Eq. (17)] for the parameters ϵ_1 and $\epsilon_2 + \epsilon_3 \equiv \epsilon'$. Experimental upper limits for $\eta\gamma$ decay modes of φ and ω permitted the rejection of sets (c) and (d), and furthermore restricted the range of ϵ_4. Obviously, measurements of these rates are highly desirable for checking the consistency of our scheme. In fact, the quoted upper limit for the rate $\Gamma(\varphi \to \eta\gamma)$ already coincides with the prediction in the SU(6) symmetry limit.

The sharpest experimental test of our scheme, however, would be the measurement of the partial decay width for $K^{*0} \to K^0\gamma$, since the possible values (a) and (b) in Eq. (18) allowed by our model are, respectively, one and two orders of magnitude lower than the symmetry limit[10] of 0.23 MeV. Measurement of $K^{*+} \to K^+\gamma$ is also important, since it determines ϵ_2 which enters in other processes, such as $K^* \to K\pi\pi$ and the K-meson electromagnetic mass shift.

Finally, further tests of the current-mixing hypothesis are provided by $\rho \to \pi\gamma$ and $\rho \to \eta\gamma$ when these are compared, respectively, with $\omega \to \pi\gamma$ and $\eta \to \pi^+\pi^-\gamma$.

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[1]M. Gell-Mann, D. Sharp, and W. G. Wagner, Phys. Rev. Letters 8, 261 (1962).

[2]See, e.g., J. Yellin, Phys. Rev. 147, 1080 (1966).

[3]C. Bemporad et al., Phys. Letters 25B, 380 (1967). A slightly improved value due to the same authors is quoted by A. H. Rosenfeld et al., Rev. Mod. Phys. 40, 77 (1968).

[4]In this work we neglect the phenomenon of X^0-η mixing. Our reasons for doing so are that (a) the mixing angle determined by the deviation from the Gell-Mann–Okubo mass formula is small and poorly determined as long as the absolute electromagnetic mass shifts are unknown; (b) attempts to explain the η lifetime by this mechanism lead to unreasonably large X^0 widths. See, e.g., A. Baracca and A. Bramon, Nuovo Cimento 51A, 873 (1967).

[5]Cf. M. Muraskin and S. L. Glashow, Phys. Rev. 132, 482 (1963).

[6]S. Coleman and H. J. Schnitzer, Phys. Rev. 134,

B863 (1964).

[7]N. M. Kroll, T. D. Lee, and B. Zumino, Phys. Rev. 157, 1376 (1967).

[8]T. D. Lee and B. Zumino, Phys. Rev. 163, 1667 (1967).

[9]The angle usually denoted by θ_Y is given in the vector-mixing model by $\tan\theta_Y = (K_\omega/K_\varphi)^{1/2}\tan\theta$, whence $\theta_Y = 33.2°$.

[10]By the SU(6)-symmetric limit we mean $K_i = 1$, $\sin\theta = 1/\sqrt{3}$, $\epsilon_i = 0$. Some of the results obtainable in this limit have been given previously by, e.g., V. V. Anisovich et al., Phys. Letters 16, 194 (1965).

[11]Rosenfeld et al., Ref. 3.

[12]The ratios of these rates are, of course, in agreement with the ratios given in R. J. Oakes and J. J. Sakurai, Phys. Rev. Letters 19, 1266 (1967).

[13]We remark that the ratio $\Gamma(K^* \to K\pi)/\Gamma(\rho \to \pi\pi)$ is also improved by the use of vector mixing. See, e.g., J. J. Sakurai, Phys. Rev. Letters 19, 803 (1967).

[14]It was already noticed that the calculation of the η lifetime using $\omega \to \pi\gamma$ [F. A. Berends and P. Singer, Phys. Letters 19, 249, 616(E) (1965)] gave a very different value from that obtained on the basis of the π^0 lifetime. The shorter η lifetime they predicted turned out to be confirmed experimentally.

[15]J. S. Lindsey and G. A. Smith, Phys. Letters 20, 93 (1966); S. M. Flatté et al., Phys. Rev. 145, 1050 (1966).

[16]We remind the reader that the calculation of rates for decays of the type $V \to P + P$ involves knowledge of $g_{VPP}(p^2)$, where $g = K_V^{1/2} g_{VPP}(0)$. Making a linear extrapolation in p^2 we obtain $g_{VPP}(p^2) = (g/\sqrt{K_V})(1-0.059p^2/m^2)$ by comparison with the well-determined K^* width. The resulting widths for ρ and φ are $\Gamma_\rho = 130 \pm 26$ MeV, $\Gamma(\varphi - \bar{K}K) = 5.1 \pm 1$ MeV.

THRESHOLD PION PRODUCTION IN NUCLEON-NUCLEON COLLISIONS

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We present results pertaining to threshold π production in nucleon-nucleon collisions, obtained in the soft-pion approximation. A typical result is $\sigma_{\rm tot}(pp \to np\pi^+) = 24\ \mu$b at $T_{\rm lab} = 310$ MeV, in good agreement with extrapolations from higher energy data.

It is possible to establish the expected[1] soft-pion connection between π production and nucleon-nucleon scattering by considering the algebra of currents as applied to the S-matrix element for the process $\gamma 2N \to \pi 2N$. Partial conservation of the axial-vector current (PCAC) is used to replace matrix elements of the pion source current by those for the axial-vector current. This facilitates the performance of certain low-energy expansions, which in lowest order, from pole terms, yield the production amplitude. The formal results are precisely consistent with those following from the Adler-Dothan theorem[2] for the weak axial-vector vertex. Calculations of the unbound process $pp \to np\pi^+$ are quite satisfactory from threshold ($T_{\rm lab} = 290$ MeV) to 320 MeV, in good agreement with the Mandelstam[3] and Rosenfeld[4] extrapolations from data existing at higher energies. The procedure can easily be extended to other processes of interest, such as $np \to pp\pi^-$ or $pp \to d\pi^+$. Our computations indicate that the situation at production threshold is compatible with results for other soft-pion processes.[5] This is in contradiction with recent work by Beder.[6]

The S-matrix element for the process $\gamma 2N \to \pi 2N$ is

$$S_{fi} = \langle \text{out}; qp_3p_4 | kp_1p_2; \text{in}\rangle \equiv \langle \text{out}; q\beta | k\alpha; \text{in}\rangle. \tag{1}$$

We suppress the charge states for the moment but later specialize to π^+ production and initial, two-proton states. The standard reduction technique using PCAC gives

$$\begin{aligned} i(2k_0 2q_0)^{\frac{1}{2}} C\langle \text{out}; \pi^\pm(q)\beta | j_\mu^{EM}(0) | \alpha; \text{in}\rangle \\ = -\int d^3z\, e^{-i\vec{q}\cdot\vec{z}}(-q^2+\mu^2)\langle \text{out}; \beta | [j_\mu^{EM}(0),\ \bar{j}_0^{(\pm)}(0,\vec{z})] | \alpha; \text{in}\rangle \\ -iq^\nu \int dz\, e^{iqz}(-q^2+\mu^2)\langle \text{out}; \beta | (j_\mu^{EM}(0)\bar{j}^{(\pm)}(z))_+ | \alpha; \text{in}\rangle, \end{aligned} \tag{2}$$

where $j_\mu^{EM} = j_\mu^3 + (\sqrt{3}/2) j_\mu^8$, $\bar{j}_\nu^{(\pm)} = \bar{j}_\nu^1 \pm i\bar{j}_\nu^2$, in terms of SU(3) vector- and axial-vector currents,