8

ИНСТИТУТ ЯДЕРНОЙ ФИЗИКИ СО АН СССР

V.M. Khatsymovsky

HYPERON RADIATIVE DECAYS IN QCD

PREPRINT 86-24

НОВОСИБИРСК

HYPERON RADIATIVE DECAYS IN QCD

V.M.Khatsymovsky

Abstract

Amplitudes of the decays $\Sigma^+ p\gamma$, $\Xi^-\Sigma^-\gamma$ and $\Omega^-\Xi^-\gamma$ are estimated in the framework of QCD sum rules method. Using this method, we are able to explain comparatively large branching ratio, $Br(\Sigma) \approx 10^{-3}$, and large negative asymmetry parameter $\alpha(\Sigma) \approx -1$ observed in the decay $\Sigma^+ p\gamma$. We also estimate branching ratios for Ξ^-, Ω^- to be (no more than) an order of magnitude smaller than the experimental upper bounds for them

1. Introduction

To describe properties of the lowest hadron states the QCD sum rules method originally proposed in ref./1/ is now widely used. Meson /1,2/ and baryon /3-5/ masses, meson formfactors and couplings /6-8/ had been calculated using this method. In refs./9,10/ the QCD sum rules for polarization operator of nucleon current in an external electromagnetic field were applied to calculation of nucleon and hyperon magnetic moments. Analogous sum rules in an external axial field were used in ref./11/ to calculate vector and axial constants of octet baryons. In ref. /12/ S-wave amplitudes for non-leptonic hyperon decays and parity violating πNN constant were calculated using sum rules for correlator of two baryon currents in presence of weak interaction ("weak field").

In this paper we derive sum rules (SR) for correlator of two baryon currents in presence of weak and electromagnetic fields. These SR allow one to calculate real parts of amplitudes of the decays $\Sigma^+ p\gamma$, $\Xi^-\Sigma^-\gamma$ and $\Omega^-\Xi^-\gamma$. These decays have the common feature in that the (lowest) pole contribution to them is small as compared to the non-pole one. The real part of the latter is just calculable with the help of SR. It turns that the dominant contribution to $\Sigma^+ p\gamma$ comes from interaction of photon with soft quarks. The largeness of $Br(\Sigma) \approx 10^{-3}$ is explained by the large value of the magnetic susceptibility of q-quark condensate, $\chi_q \approx 8$ GeV^2, introduced in /9/. Large asymmetry, $\alpha(\Sigma) \approx -1$, is explained by the large SU_3-violation due to occurence of high powers of masses m_Σ and m_p in the SR. The dominant contribution to $\Xi^-\Sigma^-\gamma$ and $\Omega^-\Xi^-\gamma$ comes from interaction of the pair of soft quarks $d, \bar{s}$ with both the weak and electromagnetic fields. Taking into account also imaginary parts of amplitudes /13-16/ we predict branching ratios for these decays.

The main idea of calculation is not new. If there are small momentum transfers in the process, corresponding distances can be large. Therefore the standard Wilson operator product expansion (OPE) /17/ for the correlator of currents must be supplemented with the vacuum expectation values (VEV's) of the nonlocal operators /18/. In our case the VEV's $\langle O \rangle_\lambda =$

$= \int d^4y [iky - \frac{1}{2}(ky)^2] \langle T\{J_\lambda(y) \mathcal{O}(0)\rangle$, $\langle \mathcal{O}\rangle_W = -i\int d^4z \cdot \langle T\{H(z)\mathcal{O}(0)\}\rangle$, $\langle \mathcal{O}\rangle_{W\lambda} = \int d^4y d^4z [iky - \frac{1}{2}(ky)^2] \cdot \langle T\{J_\lambda(y) H(z) \mathcal{O}(0)\}\rangle$ must be accounted for in addition to the usual local ones, $\langle O\rangle$. Here J_λ is the electromagnetic current, H is the four-quark $\Delta s = -1$ weak Hamiltonian.

2. The method

Our main object is the following correlator of J_λ, H and of the two three-quark baryon currents $\eta, \overline{\eta}$:

$$K_\lambda = \int d^4x d^4y d^4z \exp[i(qx + ky + \ell z)] \cdot \langle T\{\eta(x) J_\lambda(y) H(z) \overline{\eta}(0)\}\rangle \qquad (1)$$

The substitution $K_\lambda \to K_{\lambda\mu}$, $\overline{\eta} \to \overline{\eta}_\mu$ is implied in case of Ω ; ℓ is some auxiliary momentum introduced for the purpose we are going to discuss now.

At first sight, the correlator $K_{0\lambda} = K_\lambda(\ell = 0)$ seems to be more appropriate to our problem. However, it's analytic properties in variables $S_1 = -(q+k)^2$, $S_2 = -q^2$ are not quite good. The corresponding spectral density $\rho(S_1, S_2)$ is given by diagrams of the topology shown in figs. 1a, 1b. Note that such the topology corresponds to the case of $\Sigma^+ p \gamma$ only since there are no valence u -quarks in the two remaining decays. The diagrams of fig. 1b contain divergence in loop L. In terms of $K_\lambda(\ell \neq 0)$ it corresponds to the divergent dispersion integral in variable $S_\Sigma = -(q + l + k)^2$. In other words, contribution from higher states in strange channel is large. To get rid of it the Borel transformation in S_Σ (the corresponding parameter is denoted as Σ) is suitable /1/. Putting $l^2 = kl = k^2 = 0$ we are left generally speaking with the only variable $s_p = -q^2 = -(q+k)^2$ (the matter is that the contribution of J -induced VEV $\langle O\rangle_\lambda$ dominates and is calculable in kinematic $k = 0$ only). Then borelization in s_p suppresses continuum in the proton channel.

However, to find the phisical amplitude A of interest from the SR one needs some model assumptions concerning the dependence of A on S_p, S_Σ. In the aspect of bordization in S_p, S_Σ the behaviour of A in the vicinity of $S_p = -m_p^2$, $S_\Sigma = -m_\Sigma^2$ is of interest. So we represent A approximately as the sum of the pole contribution and some constant term A_0: $A = A_0 + (S_p + m_p^2)^{-1} A_p + (S_\Sigma + m_\Sigma^2)^{-1} A_\Sigma$. Here A_p, A_Σ are known from the pole graphs of fig.2, and A_0 is the object of our calculation. After transferring the continuum to the RHS and borelization we have the LHS, phenomenological part, of the SR to be proportional to

$$\frac{1}{P}\left(A_0 + \frac{A_P}{P} + \frac{A_\Sigma}{\Sigma}\right) \frac{1}{\Sigma} \exp\left(-\frac{m_P^2}{P} - \frac{m_\Sigma^2}{\Sigma}\right) \qquad (2)$$

In what follows we compare our SR with the SR for some known physical quantities and reduce them to the form $A_0 + A_P P^{-1} + A_\Sigma \Sigma^{-1}$ = const + higher power corrections. The terms $\sim A_P, A_\Sigma$ are not essential for they can be either attributed to the higher power corrections here not accounted for or cancelled by applying the operator ($\partial^2/\partial P \partial \Sigma) \cdot (P\Sigma \cdot$) to both sides of the SR. In the duality estimates $P, \Sigma \to \infty$ and these terms again drop out. However, the pole contribution is to be added to A_0 to give the physical amplitude in the Minkovsky region.

3. Nonlocal VEV's

Consider first the correlators in the weak field. VEV's $\langle O\rangle_W$ can be calculated using equations of motion, $i\not{D}q = m_q q + H_{\overline{q}}$, $H_{\overline{q}} \equiv \partial H/\partial \overline{q}$. For example, $0 \equiv i\partial^\mu \langle \overline{s}\gamma_\mu d\rangle_W = (m_d - m_s)\langle \overline{s}d\rangle_W + \langle(\overline{s}H_{\overline{d}} - H_s d)\rangle$ which gives

$$\langle \overline{s}d\rangle_W = \left\langle \frac{\overline{s}H_{\overline{d}} - H_s d}{m_s - m_d} \right\rangle \qquad (3)$$

The RHS of (3) is the VEV of some local four-quark operator $\mathcal{O}_4$. At the normalization momenta $M \sim 1 GeV \sim m_c$ this operator contains only left-handed quark fields because H does so 19 (M is the typical scale of Borel parameters, m_c is the C -quark mass). Therefore being factorized at M (3) vanishes. However, dressing $\mathcal{O}_4$ with gluons of momenta from m_c to $\mu < m_c$ leads to new operators of the kind $\overline{\psi}_L \psi_R \overline{\psi}_R \psi_L$ with the typical coefficient $C_R = \frac{1}{6\pi} \ln m_c/\mu$ /20/. Therefore factorization $\langle \mathcal{O}\rangle_4$ at μ gives already nonzero result.

Likely VEV $\langle \nabla_\mu d\overline{s}\rangle_{W\lambda}$ which gives the dominant contribution to Ω^- and Ξ^- decays can be reduced to $\langle \overline{s}_L \sigma_{\alpha\beta} H_{\overline{d}}\rangle_\lambda$ and $\langle H_s \sigma_{\alpha\beta} d_L\rangle_\lambda$ ($\sigma_{\alpha\beta} = \frac{1}{2}[\gamma_\alpha, \gamma_\beta]$). The latter can then be factorized at $\mu < m_c$: $\langle \psi_L \overline{\psi}_R \psi_R \overline{\psi}_L \rangle_\lambda = \langle \psi_L \overline{\psi}_R \rangle_\lambda \langle \psi_R \overline{\psi}_L \rangle + \langle \psi_L \overline{\psi}_R \rangle \langle \psi_R \overline{\psi}_L \rangle_\lambda$. But in case of $\langle \bar{s}\sigma d \rangle_{w\lambda}$ the equations of motion are useless. Consider this correlator at $k^2 \sim -1$ GeV^2. OPE for it begins with $\langle \bar{s} d \rangle_w$. Analogous OPE for $\langle \bar{q}\sigma q \rangle$ /21/ begins with $\langle \bar{q} q \rangle$. Comparing them gives

$$\langle \bar{s}\sigma d \rangle_{w\lambda} \approx \langle \bar{d}\sigma d \rangle_\lambda \langle \bar{s} d \rangle_w \langle \bar{d} d \rangle^{-1} \tag{4}$$

up to the power corrections not accounted for. In the framework of the lowest vector state dominance one can then extrapolate (4) to $k^2 = 0$ with an accuracy of 1.5.

The linear in k terms in $\langle O \rangle_\lambda$ were obtained in /9,10/. In analysis of Ω^- decay the bilinear in k terms in $\langle O \rangle_\lambda$, $\langle O \rangle_{w\lambda}$ are also of interest. Using C-conjugation we obtain, for example, that $\langle \bar{u}\sigma_{\alpha\beta} \nabla_\mu u \rangle_\lambda = \langle \overline{\nabla_\mu u}\, \sigma_{\alpha\beta} u \rangle_\lambda = \frac{1}{2} \langle \partial_\mu (\bar{u}\sigma_{\alpha\beta} u) \rangle_\lambda = \frac{1}{2} i k_\mu \langle \bar{u}\sigma_{\alpha\beta} u \rangle_\lambda$ and so on. Some formulas for the VEV's considered are presented in Appendix A. Multi-quark VEV's are supposed to obey factorization hypothesis, for instance

$$\langle d\bar{s} u\bar{u} \rangle_w = \langle d\bar{s} \rangle_w \langle u\bar{u} \rangle \tag{5}$$

Here the fact is taken into account that the contribution to the weak interaction of four valence quarks d, $\bar{s}$, u, $\bar{u}$ at large distances z is suppressed by z^{-6} as compared to that of two valence quarks d, $\bar{s}$. The former becomes significant in higher orders of OPE /12/.

Inclusion of the effects of H at large distances into SR is straightforward for Ω^- and Ξ^- decays for which the short distances are unessential. Here we work in kinematic l=0 from the very beginning, since H- and (H + J)-induced VEV's are known in this kinematic only. As for the Σ^+ decay, we need to consider it to consist of two subprocesses caused by H at different distances. The diagrams dominating in each case differ in their topology, and this difference keeps in each order of OPE as soon as we do not exceed the limits of factorization of the VEV's (see (5)).

4. Derivation of the sum rules

Consider phenomenological expression for K_λ and choice of γ-matrix structure needed to formulate the SR. Parametrization of the matrix element sought for, $\int d^4y \exp(iky) \cdot \langle B_2 | T\{ J_\lambda(y) H(0) \} | B_1 \rangle$, is usually chosen as

$$i G_F m_\pi \bar{u}_2 (a\gamma_5 + b) \sigma_{\lambda\rho} u_1 k^\rho \tag{6}$$

for decays of $1/2^+$ and

$$i G_F m_\pi \bar{u} [a_1 \sigma_{\lambda\rho} \gamma_5 u_\mu k^\mu + a_2 (\gamma_\lambda \gamma_5 u_\rho - \gamma_\rho \gamma_5 u_\lambda) + b_1 \sigma_{\lambda\rho} u_\mu k^\mu + b_2 (\gamma_\lambda u_\rho - \gamma_\rho u_\lambda)] \tag{7}$$

for $\Omega^- \Xi^- \gamma$. In what follows, calculations of the amplitudes with exception of $a(\Sigma)$ (it vanishes in SU_3-limit /22/) are made neglecting baryon mass difference: $m_1 = m_2 = m$. The residues of baryons into currents used are defined according to formulas $\langle 0 | \eta | \frac{1}{2}^+ \rangle = \beta \gamma_5 u$, $\langle 0 | \eta_\mu | \frac{3}{2}^+ \rangle = \lambda u_\mu$, where $\bar{u}u = 2m$, $\bar{u}_\mu u_\mu = -2m$. To calculate contribution of H at short distances (it is the case of $\Sigma^+ p \gamma$) we take the structures $q_\lambda [\not{q}, \not{k}] \left\{ \begin{matrix} 1 \\ \gamma_5 \end{matrix} \right.$ having the maximal number of momenta. The dominating contribution to corresponding SR is due to $\langle \psi\overline{\psi} \rangle_\lambda$. It is given by the two-loop diagrams of the type shown in fig.3c. The SR for the structures next in number of momenta, $\not{q}\gamma_\lambda \not{k} - \not{k}\gamma_\lambda \not{q}$ or $q_\lambda \not{k} \gamma_5$ begin with the asymptotic loop of fig.3a which is considerable equally with $\langle \psi\overline{\psi}\psi\overline{\psi} \rangle_\lambda$- terms (see fig.3b) and very sensitive to the choice of the model for continuum due to high power of M^2. Therefore the given SR are of little use for calculation, though explicit estimates are indicative of consistency of these SR with the former ones up to the factor of 2.

The weak interaction at large distances leads to the diagrams having no more than one loop (see figs. 3g, 3h, 3i, 3j). Therefore the continuum is not of a great value here. Let as chose the structures so as to suppress maximally contribution to the SR from transition between baryons of opposite P-parity (this contribution is not sufficiently suppressed by ordinary borelization). Our aim is achieved if the amplitude of interest enters coefficient at the chosen structure being multiplied by the sum of baryon masses. The amplitude of undesirable transition $\frac{1}{2}\left(\frac{3}{2}\right)^{\pm} \to \frac{1}{2}^{\mp}$ is then multiplied by the difference of masses. Such the property is characteristic of the structures $\not{q}\gamma_\lambda \not{k} - \not{k}\gamma_\lambda \not{q}$ and $q_\lambda \not{k}\gamma_5$ in K_λ and of the following five structures in $K_{\lambda\mu}$ here borelized in S_1, S_2:

$$K_{\lambda\mu} = iG_F \frac{\beta\lambda}{t_1 t_2} exp\left(\frac{m^2}{t}\right)[(-2ma_1 + a_2)k_\mu q_\lambda \not{k} - \tag{8}$$
$$-\frac{1}{3} m a_2 (k_\mu[\gamma_\lambda, \not{q}] + g_{\lambda\mu}[\not{q}, \not{k}]) - (mb_1 + \frac{1}{2} b_2)\cdot$$
$$\cdot k_\mu (\not{q}\gamma_\lambda \not{k} - \not{k}\gamma_\lambda \not{q})\gamma_5 + \frac{4}{3} m b_2 k_\mu q_\lambda \gamma_5 + \ldots]$$

Here t_1, t_2 are the Borel parameters, $t = t_1 t_2 (t_1 + t_2)^{-1}$. Further, we can perform at $l = 0$, generally speaking, only the ordinary Borel transformation in $s = s_1 = s_2$ (when contribution of $\langle \mathcal{O} \rangle_{W\lambda}$ is calculated). However, to take into account the continuum contribution more accurately, we shall present the contribution of $\langle \mathcal{O} \rangle_W$ in the twice borelized in s_1, s_2 form. As for the single borelized at $l = k = 0$ SR, taking into account continuum in them consists in the usual substitution $t^n \to$ $\to [(n-1)!]^{-1} \int exp(-s/t) s^{n-1} \theta[s(s_0 - s)] ds$ in the RHS on condition that there is $A\ exp(-m^2/t)$ in the LHS (t is the Borel parameter). Such the model for continuum corresponds to the certain form of the double spectral density, $\rho(s_1, s_2) = \delta(s_1 - s_2) f(s_1)$. Then the double Borel transformation in s_1, s_2 is equivalent to the single one in $s = s_1 = s_2$ with the subsequent substitution $t \to t_1 t_2 (t_1 + t_2)^{-1}$. This substitution is just implied below.

Let us denote by sw and ℓw contributions to the amplitudes of $\Sigma^+ p \gamma$ from the weak interaction at short and at large distances, correspondingly. Besides that, $\tilde{\beta} = (2\pi)^2 \beta$, $\tilde{\lambda} = (2\pi)^2 \lambda$, $a_q = -(2\pi)^2 \langle \bar{q} q \rangle$, $m_0^2 = i g_s \langle \bar{q} \sigma G q \rangle \langle \bar{q} q \rangle^{-1}$, χ_q is the q-quark condensate magnetic sucseptibility /9/, $C_-(M) = (\ln \frac{m_W}{\Lambda} / \ln \frac{M}{\Lambda})^{4/9}$. Calculating a_{sw}, b_{sw} we have found contribution to them from VEV's $\langle \Psi\bar{\Psi} \rangle_\lambda$, $\langle \Psi\bar{\Psi} \rangle$, $\langle \Psi\bar{\Psi} G \rangle_\lambda$ and from $\langle \Psi\bar{\Psi}\Psi\bar{\Psi}\Psi\bar{\Psi} \rangle_\lambda$ singled out by the absence of loop smallness (see figs. 3e, 3d, 3e, 3f correspondingly):

$$b_{sw} - a_{sw} = -\langle \bar{d} d \rangle \frac{cs\sqrt{2}\, C_-(M)}{18 m_\pi \tilde{\beta}_\Sigma \tilde{\beta}_p} exp\left(\frac{m_\Sigma^2}{\Sigma} + \frac{m_p^2}{p}\right)\cdot \tag{9}$$
$$\cdot\left[\chi_d \frac{\Sigma^3 p}{\Sigma + p} + \Sigma^3 \frac{\Sigma - p}{(\Sigma + p)^2} + 4\Sigma^2 - \left(\frac{\Sigma p}{\Sigma + p}\right)^2 + \right.$$

$$\left. + \frac{1}{6} m_{0d}^2 \chi_d \left(\Sigma^2 - \frac{1}{2}\frac{\Sigma^3}{\Sigma + p}\right) + \frac{4}{3} a_u^2 (2\chi_u + \chi_d)\right]$$

To derive the SR for $a_{\ell w}, b_{\ell w}, a(\Xi), b(\Xi)$ and a_1, b_1, a_2, b_2 we have calculated contribution from VEV's $\langle d(x)\bar{s}(0) \rangle_W$, $\langle d(x)\bar{s}(0)\Psi(0)\bar{\Psi}(0) \rangle_W$, $\langle d(x)\bar{s}(0) \rangle_{W\lambda}$, $\langle d(x)\bar{s}(0)\Psi(0)\bar{\Psi}(0) \rangle_{W\lambda}$ (see diagrams of figs. 3g, 3h, 3i, 3j respectively). Of these only $\langle \nabla_\mu d\bar{s} \rangle_W$ contributes to $\Sigma^+ p \gamma$:

$$\left.\begin{matrix} a_{\ell w} \\ b_{\ell w} \end{matrix}\right\} = \frac{8}{27} \frac{C_R}{\pi^2} \frac{a^2 cs\sqrt{2}}{m\tilde{\beta}^2} t_1 t_2 \frac{t_2 \mp t_1}{(t_2 + t_1)^2} exp\left(\frac{m^2}{t}\right) \tag{10}$$

SR for Ξ^-, Ω^- are given in Appendix B.

5. Estimates

Expression (9) for $b_{sw} - a_{sw}$ does not contain any positive power of p when p is large. This indicates the smallness of the continuum contribution in the proton channel. Therefore we can put $p = \infty$ in this SR and, analogously, $\Sigma = \infty$ in $b_{sw} + a_{sw}$. Further, let us compare our SR with the baryon mass SR /4/,

$$\begin{aligned} 2\tilde{\beta}^2 &= (t^3 + 4a^2/3)\, exp(m^2/t) \\ m\tilde{\beta}^2 &= at^2\, exp(m^2/t) \\ 3m\tilde{\lambda}^2 &\approx 4at^2\, exp(m^2/t) \end{aligned} \tag{11}$$

and with the hyperon s-wave SR /12/,

$$A(\Sigma_0^+ + \Lambda_-^0 \sqrt{3}) \approx \frac{4}{3} \frac{(2\pi)^2}{m_\pi^2 f_\pi} \frac{\langle \bar{s} d \rangle_W\, at}{G_F \sqrt{2}\, m\tilde{\beta}^2} exp\left(\frac{m^2}{t}\right) \tag{12}$$

This gives

$$b_{sw} = \frac{cs\sqrt{2}\, C_-(\sqrt{s_0})}{(6\pi)^2 m_\pi} a\chi \left[1 + \frac{3m}{2a\chi} + \frac{m_0^2}{12}\frac{m}{a} + \right.$$
$$\left. + \frac{4}{3}\frac{a^2}{\tilde{\beta}^2}\right] = 7.2 \cdot 10^{-2} \tag{13}$$
$$b_{\ell w} = A(\Sigma_0^+ + \Lambda_-^0 \sqrt{3}) \frac{m_\pi f_\pi}{a\sqrt{2}} \frac{m_s}{f} = 2.1 \cdot 10^{-2}$$

$$b(\Xi)=\frac{2\sqrt{2}}{9}csC_R\frac{a\chi}{\pi^2 m_\pi}=\frac{2a_2}{\sqrt{3}}=5.4\cdot 10^{-2}$$

$$a(\Xi):b(\Xi)\simeq 1:10,\ a_1:a_2:b_1:b_2\simeq 3GeV^{-1}:6:1GeV^{-1}:(-1)$$

where the numerical estimates are presented for $\chi = 8GeV^{-2}$, /9/, $C_R=0.25$ /18/, $m_0^2=0.8GeV^2$ /4/, $a=0.55GeV^3$, $m_s=0.15GeV$, $f=\langle(\bar{s}s-\bar{d}d)\rangle\langle\bar{d}d\rangle^{-1}=-0.2$, $C_-(\sqrt{s_0})=1.5$ (at $m_W=80GeV$, $\sqrt{s_0}=1.5GeV$, $\Lambda=0.1GeV$); $\tilde{\beta}^2 = 1.1GeV^6$ is obtained using the SR of ref. /4/ at m = 0.94 GeV. Besides that, the pole contribution to $b(\Sigma)$ should be also taken into account (see section 2). It is estimated to be $b_{pole} = -2.4\cdot 10^{-2}$. To arrive at this number, we have used the experimental values of $\langle p|H|\Sigma^+\rangle$ known from s-waves and those of magnetic moments of p and Σ^+ substituting them into graphs of fig.2. Analogous graphs in $\Xi^-\Sigma^-\gamma$ and $\Omega^-\Xi^-\gamma$ are unimportant /16/.

To estimate a_{sw} the SU_3-violation must be accounted for which enters SR (9) mainly through the difference of the scales $\Sigma\simeq m_\Sigma^2$ and $p\simeq m_p^2$ or through that of the continuum thresholds $(s_{10}-s_{20})/s_{20}\simeq(m_\Sigma^2-m_p^2)m_p^{-2}$. So we get $-a_{sw}/b_{sw}\simeq +0.6$. Accounting for SU_3 -violating VEV's of the type $\langle(s\bar{s}-d\bar{d})\rangle$ leads to slight correction of order of $\langle(\bar{d}d-\bar{s}s)\rangle\cdot\langle(\bar{d}d+\bar{s}s)\rangle^{-1}\simeq +0.1$ to this value. Besides that

$$-a_{ew}/b_{ew}\simeq +0.25$$

Collecting the obtained contributions we get for the amplitudes of $\Sigma^+p\gamma$: $b(\Sigma)=b_{sw}+b_{ew}+b_{pole}=6.9\cdot 10^{-2}$, $a(\Sigma)=a_{sw}+a_{ew}=-5.5\cdot 10^{-2}$. Asymmetry parameter of $\Sigma^+p\gamma$ and branching ratios are then

$$\begin{aligned}&\alpha(\Sigma)=2ab(a^2+b^2)^{-1}=-1.0\\ &Br(\Sigma)=0.8\cdot 10^{-3}\\ &Br(\Xi)\geqslant 1.0\cdot 10^{-4}\\ &Br(\Omega)\geqslant 2.2\cdot 10^{-4}\end{aligned}\tag{14}$$

Combining these with known imaginary parts of Ξ^- - and Ω^--amplitudes /13-16/ we finally obtain

$$\begin{aligned}&Br(\Xi)=2.0\cdot 10^{-4}\\ &Br(\Omega)=2.3\cdot 10^{-4}\end{aligned}\tag{15}$$

Experimental values are as follows /23/:

$$\begin{aligned}&-\alpha(\Sigma)=0.72\pm 0.29\\ &Br(\Sigma)=(1.20\pm 0.13)\cdot 10^{-3}\\ &Br(\Xi)<1.2\cdot 10^{-3}\\ &Br(\Omega)<3.1\cdot 10^{-3}\end{aligned}\tag{16}$$

6. Conclusion

The accuracy of our calculations is governed by a number of factors. The principial feature of the situation is that from the QCD viewpoint the calculated value is the amplitude of scattering baryon by two currents, J_λ and H. Such the amplitude depends nontrivially on the internal s-t-channel variables some of which lying in the Minkovsky region. Consequently, the SR which are valid in the Enclidean region must be supplemented with some model assumptions allowing one to extrapolate the amplitude from Enclidean to Minkovsky region (the simple model of such kind used by us represents the amplitude as the sum of a constant and lowest poles). Besides that, there are the problems of calculation and renormalization of complex VEV's and those connected with uncertainty in the model for continuum in the SR at small momentum transfers. As a result, we expect for amplitudes to be determined with an accuracy up to the factor of 1.5 or even 2 in the case of Ξ^-, Ω^-.

In conclusion, the author is indebted to I.B.Khriplovich who has drawn the author's attention to the given problem, for consideration of the work; to A.R.Zhitnitsky and B.L.Ioffe for useful discussions.

Appendix A.

The terms of interest in the expansion of some nonlocal VEV's in x,k take the form:

$$\langle u_i^a(x)\bar{u}_k^b(0)\rangle_\lambda = i\frac{\langle\bar{u}u\rangle}{18}k^\rho\Big[\chi_u\sigma_{\lambda\rho}\Big(1+\frac{i}{2}kx\Big)+\frac{1}{8}\Big(1+\frac{1}{6}\chi_u m_{0u}^2\Big)\cdot x^2\sigma_{\lambda\rho}-\frac{1}{24}\not{x}\sigma_{\lambda\rho}\not{x}\Big]_{ik}\delta^{ab} \tag{1A}$$

$$\langle ig_s G^n_{\alpha\beta}u_i^a\bar{u}_k^b\rangle_\lambda = -i\frac{\langle\bar{u}u\rangle}{18\cdot 32}\chi_u m_{0u}^2 k^\rho[\{\sigma_{\lambda\rho},\sigma_{\alpha\beta}\}_+]_{ik}(t^n)^{ab} \tag{2A}$$

$$\langle d_i^a(x)\bar{s}_k^b(0)\rangle_W = \frac{4}{27}G_F cs\sqrt{2}\,C_R\langle\bar{q}q\rangle^2\Big[-f\frac{1+\gamma_5}{m_s}+\frac{i}{4}\not{x}(1-\gamma_5)\Big]_{ik}\delta^{ab} \tag{3A}$$

$$\langle d_i^a(x)\bar{s}_k^b(0)\rangle_{W\lambda} = i\frac{4}{81}G_F cs\sqrt{2}\,C_R\langle\bar{q}q\rangle^2\chi k^\rho\cdot\Big[-\frac{f}{m_s}\cdot\Big(1+\frac{i}{2}kx\Big)\sigma_{\lambda\rho}(1+\gamma_5)+\frac{3}{8}i\{\not{x},\sigma_{\lambda\rho}\}_+(1-\gamma_5)\Big]_{ik}\delta^{ab} \tag{4A}$$

Notations are given in sect.4,5.

Appendix B.

The SR for Ξ^- and Ω^- decay amplitudes take the form (four terms enumerated in brackets are given by diagrams of figs. 3g, 3h, 3i, 3j, respectively):

$$a(\Xi)=2D\Big(\frac{1}{2}t_1t_2\frac{t_1-t_2}{(t_1+t_2)^2}+0+0+0\Big) \tag{1B}$$

$$b(\Xi)=2D\Big(0+\frac{a}{3}\frac{f}{m_s}+\frac{3}{4}\chi t^2-\frac{2}{3}\chi a\frac{f}{m_s}\cdot t\Big) \tag{2B}$$

$$a_2=\frac{1}{2}D\Big(-3\frac{t_1t_2}{t_1+t_2}+a\frac{f}{m_s}+3\chi t^2+0\Big) \tag{3B}$$

$$b_2=\frac{3}{4}D\Big(t_1t_2\frac{t_1-t_2}{(t_1+t_2)^2}-a\frac{f}{m_s}+0+\frac{2}{3}\chi a\frac{f}{m_s}t\Big) \tag{4B}$$

$$a_1-\frac{a_2}{2m}=D\Big(-\frac{2f}{3m_s}t_1t_2^2\frac{3t_1-t_2}{(t_1+t_2)^3}-\frac{a}{3t_1}-\frac{\chi}{6}\frac{f}{m_s}t^2-\frac{\chi}{6}a\Big) \tag{5B}$$

$$b_1+\frac{b_2}{2m}=D\Big(0-\frac{a}{3t_1}-\frac{\chi}{6}\frac{f}{m_s}t^2-\frac{\chi}{6}a\Big) \tag{6B}$$

Here $$D=\frac{4}{27}\frac{C_R}{\pi^2}\frac{a^2cs\sqrt{2}}{m\tilde{\beta}^2}\exp\Big(\frac{m^2}{t}\Big).$$

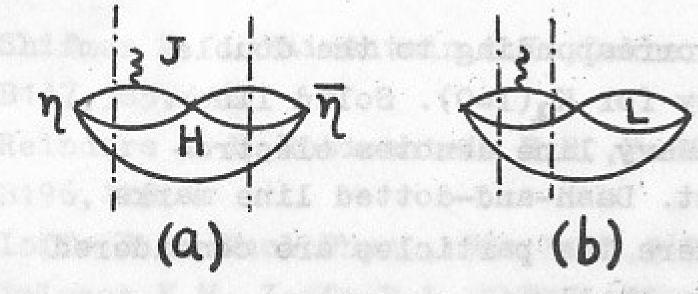

Fig.1.

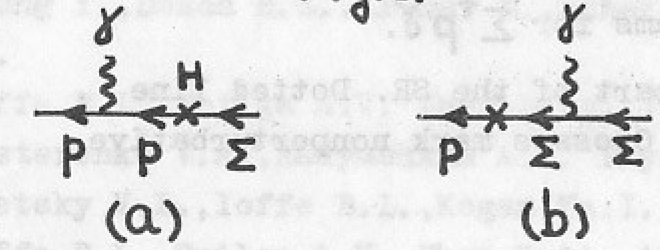

Fig.2.

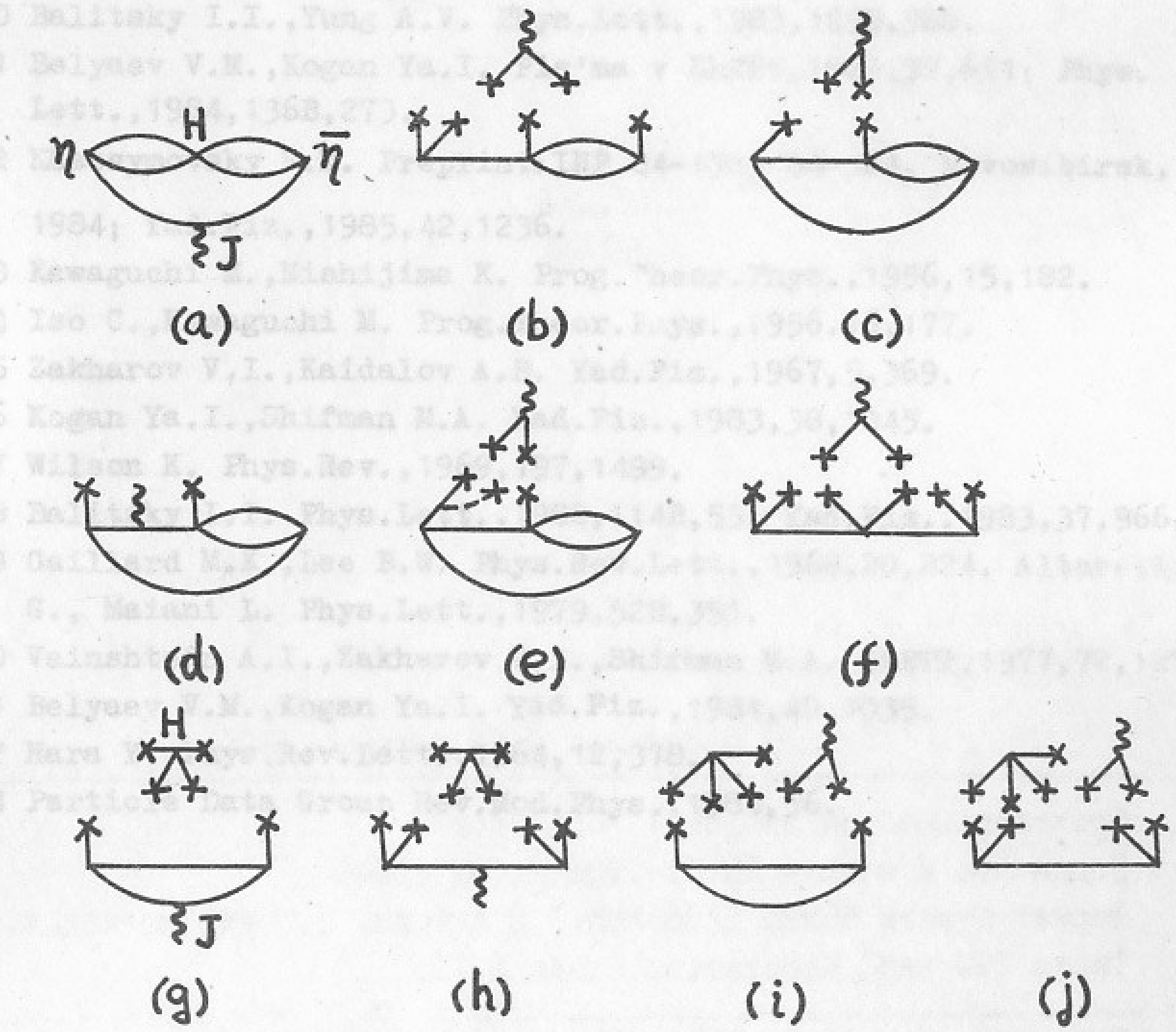

Fig.3.

Figure captions.

Fig.1. The diagrams, corresponding to the double spectral density for $K_\lambda(l=0)$. Solid line denotes quark. Wavy line denotes electromagnetic current. Dash-and-dotted line marks the channels where the particles are considered as lying on mass shell.

Fig.2. The pole diagrams for $\Sigma^+ p \gamma$.

Fig.3. "Theoretical" part of the SR. Dotted line denotes gluon. Crosses mark nonperturbative fields.

References

1 Shifman M.A.,Vainshtein A.I.,Zakharov V.I. Nucl.Phys.,1979, B147,385,448.

2 Reinders L.J.,Rubinshtein H.R.,Yazaki S. Nucl.Phys.,1982, B196,125.

3 Ioffe B.L. Nucl.Phys.,1981,B188,317; (E),1981,B191,591.

4 Belyaev V.M.,Ioffe B.L. ZhETF,1982,83,876; ZhETF,1983,84,1236.

5 Chung Y.,Dosch H.G.,Kremer M.,Schall D. Nucl.Phys.,1982,B197, 55.

6 Ioffe B.L.,Smilga A.V. Phys.Lett.,1982,114B,353.

7 Nesterenko V.A.,Radyushkin A.V. Phys.Lett.,1982,115B,410.

8 Eletsky V.L.,Ioffe B.L.,Kogan Ya.I. Phys.Lett.,1983,122B,423.

9 Ioffe B.L.,Smilga A.V. Phys.Lett.,1983,133B,436 Nucl.Phys., 1984,B232,85.

10 Balitsky I.I.,Yung A.V. Phys.Lett.,1983,129B,388.

11 Belyaev V.M.,Kogan Ya.I. Pis'ma v ZhETF,1983,37,611; Phys. Lett.,1984,136B,273.

12 Khatsymovsky V.M. Preprints INP 84-138, 84-164, Novosibirsk, 1984; Yad.Fiz.,1985,42,1236.

13 Kawaguchi M.,Nishijima K. Prog.Theor.Phys.,1956,15,182.

14 Iso C.,Kawaguchi M. Prog.Theor.Phys.,1956,16,177.

15 Zakharov V.I.,Kaidalov A.B. Yad.Fiz.,1967,5,369.

16 Kogan Ya.I.,Shifman M.A. Yad.Fiz.,1983,38,1045.

17 Wilson K. Phys.Rev.,1969,197,1499.

18 Balitsky I.I. Phys.Lett.,1982,114B,53; Yad.Fiz.,1983,37,966.

19 Gaillard M.K.,Lee B.W. Phys.Rev.Lett.,1968,20,224. Altarelli G., Maiani L. Phys.Lett.,1979,52B,351.

20 Vainshtein A.I.,Zakharov V.I.,Shifman M.A. ZhETF,1977,72,1275.

21 Belyaev V.M.,Kogan Ya.I. Yad.Fiz.,1984,40,1035.

22 Hara Y. Phys.Rev.Lett.,1964,12,378.

23 Particle Data Group Rev.Mod.Phys.,1984,56.

В.М.Хацимовский

РАДИАЦИОННЫЕ РАСПАДЫ ГИПЕРОНОВ В КХД

Препринт
№ 86-24

Работа поступила - 23 декабря I985 г.

Ответственный за выпуск - С.Г.Попов
Подписано к печати 28.0I-I986 г. МН II644
Формат бумаги 60х90 I/I6 Усл.I,3 печ.л., I,0 учетно-изд.л.
Тираж 290 экз. Бесплатно. Заказ № 24.

Ротапринт ИЯФ СО АН СССР, г.Новосибирск, 90