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Abstract

S-wave amplitudes of the non-leptonic hyperon decays are calculated in the soft pion limit in a modelless way using the QCD sum rules method. The results obtained agree with the experimental ones with an accuracy $\lesssim 20\%$.

1. Introduction

Recently the QCD sum rules method suggested in the paper [1] is widely used to describe properties of the lowest hadron states. The masses of mesons [1-2] and of baryons [3-5], form-factors and meson coupling constants [6-8] have been calculated. In refs. [9,10] the QCD sum rules for polarization operator of nucleon current in an external electromagnetic field first suggested in ref. [9] were used for calculation of the magnetic moments of octet baryons. In ref. [11] using analogous sum rules in an external axial field the axial constants of octet baryons were calculated.

In this paper the QCD sum rules method is used to determine the matrix elements of the weak Hamiltonian between the octet baryons, $\langle B_2|H_W|B_1\rangle$. Knowing these matrix elements one knows according to PCAC the S-waves in the soft pion limit. The values for S-waves obtained are consistent with the experimental ones with an accuracy of 20%.

Acting in spirit of QCD sum rules to analyse a weak transition $B_1 \to B_2$ one would have to consider the Wilson operator expansion for the T-product of the two baryon currents. For example, if the transition $\Sigma^+ \to p$ is considered, the following correlator is of interest

$$K = i\int dx\, e^{iqx}\langle 0|T\{\eta_p(x)\bar{\eta}_\Sigma(0)\}|0\rangle = \sum_n \langle 0|\mathcal{O}_n|0\rangle (C_n \hat{q} + D_n) \qquad (1)$$

where $\mathcal{O}_n$ are local operators, $\eta_p, \bar{\eta}_\Sigma$ are operators with the proton and Σ-hyperon quantum numbers, respectively. The RHS of (1) is decomposed in the two independent γ-matrix structures 1, $\hat{q}$ and their coefficients C_n, D_n are power functions of $s = -q^2$. It is clear that the polarization operator (1) differs from zero only if the weak interaction is taken into account, and we are interested in the first order of its expansion in the powers of Fermi constant G_F. The dependence on G_F is contained in both the vacuum averages $\langle 0|\mathcal{O}_n|0\rangle$ and in the coefficient functions C_n, D_n. An alternative expression for K is provided by the general dispersion relation which gives the correlator in terms of phenomenological matrix elements $\langle p_i|H_w|\Sigma_i\rangle$, $\langle 0|\eta_p|p_i\rangle$, $\langle\Sigma_i|\bar{\eta}_\Sigma|0\rangle$ where Σ_i, p_i are the real states created by currents $\bar{\eta}_p$, η_Σ from vacuum. Equating both the expressions for K one arrives at QCD sum rules. For definiteness we shall refer to the phenomenological part of sum rules as to their LHS and by RHS we mean the power expansion of the correlator.

So formulated the problem would be analogous to that of calculation of baryon magnetic moments [9,10] and axial constants [11] by the external field method. Now some scalar-pseudoscalar field coupled to H_w plays the role of an external one. However, some complications arise in our case due to occurence of nonpolinomial in $s = -q^2$ subtraction terms in the RHS of the sum rules. These terms are to be added, for example, to the diagram of fig. 1 proportional to

$$\propto s^3 \ln^2 \frac{\Lambda_{UV}^2}{s} \qquad (2)$$

and to that of fig. 2 containing the terms like

$$\propto \frac{\Lambda_{UV}^2}{s} \qquad (3)$$

where Λ_{UV} is an ultra-violet cut-off. As a result, the dependence on Λ_{UV} can be included in a definition of the operators η_Σ, η_p at some momentum s with taking into account weak interactions. Then Λ_{UV} is substituted by some constants of order of a typical hadronic mass. It is important that physics must not depend on these constants as well as on the subtraction terms themselves. Really, these terms are given by the diagrams where the weak vertex and one of the current vertices are tightened to the same point and correspond to the case when only one of the baryon currents creates or annihilates real hadronic states (the so-called single pole terms). The commonly used Laplace transformation of the sum rules first suggested in ref. [1] is not sufficient to cancel single pole terms. Remind that the Laplace transformation of some analytical function $f(s)$ is defined as

$$f(t) \equiv L_t f(s) = \int_{\varepsilon - i\infty}^{\varepsilon + i\infty} \frac{ds}{2\pi i t} f(s) \exp\frac{s}{t}, \; \varepsilon > 0 \qquad (4)$$

It allows one 1) to improve convergence of s^{-1} series, 2) to suppress the higher states contribution in the correlator and 3) to cancel the subtraction terms polinomial in s. The only possibility to separate single pole terms is to write down the double dispersion relation for correlator considering the momentum variables $s_1 = -q_1^2$ and $s_2 = -q_2^2$ of the initial and final baryons, respectively, as independent ones. Then the double Laplace transformation in s_1, s_2 first introduced in ref. [6] turns out to be adequate to the problem. To make distinction between s_1 and s_2 a momentum $k = q_2 - q_1 \neq 0$ is introduced into the weak vertex. Then the value calculated becomes a function of k^2 but we put $k^2 = 0$ so the object of calculation is the former matrix element of interest.

Thus, the external field becomes a variable one. The procedure of the double Laplace transformation $L_1 L_2$ which is of interest for us is rather straightforward in the diagrams like those depicted in figs. 1,2 including the weak vertex at short distances. It is convenient to use the Feynman parametrization of a diagram to perform this transformation. As for the graphs relevant to a G_F-dependence of the vacuum expectation values (VEV's) $\langle 0|\mathcal{O}_n|0\rangle$ (fig. 7, see further) we can unambiguously compute them only at $k = 0$; at $k \neq 0$ we have instead of these VEV's some complicated nonperturbative vertices of an interaction of soft fields with the variable external one. However, it is possible to use the results of calculation of these graphs in the constant external field extending it in proper way to the region of noncoincident s_1, s_2. The hypothesis we use is that the contribution into the double spectral density of a diagram relevant to the interaction with the external field at large distances takes the form

$$\rho(s_1, s_2) = f(s_1)\,\delta(s_2 - s_1) \qquad (5)$$

Therefore the double Laplace transformation $L_1 L_2$ at the coincident parameters $t_1 = t_2 = 2t$ reduces to the ordinary one L_t

in $s = s_1 = s_2$ with the parameter t. Some argumentation in favour of (5) is given below. In particular, we shall see that this form of $\rho(s_1, s_2)$ (and only this form!) allows one to reproduce in the one-dimensional language $(s_1 = s_2)$ the earlier obtained sum rules [9,11] for the polarization operators in the constant external fields. Thus, (5) is no more than a formal ground of that one can calculate the contribution of the two types of diagrams into the physical amplitude using two different techniques, namely, the double and the ordinary Laplace transformations.

The paper is organized as follows. In the next section some notations and definitions are introduced. In sect. 3 we calculate the graphs which include the weak interaction at small distances. In sect. 4 we give some phenomenological arguments in favour of (5). In sect. 5 we study the weak interaction at large distances and present the final results.

2. Notations and definitions

We stick to the following definition of the S-wave amplitude A for a decay $B_1 \to B_2 \pi$

$$\langle B_2 \pi | i H_W^{PV} | B_1 \rangle = G_F m_\pi^2 A \bar{u}_2 u_1 \tag{6}$$

Then according to PCAC in the soft pion limit

$$A(\Sigma_0^+) = -\frac{\alpha}{G_F m_\pi^2 f_\pi \sqrt{2}} \tag{7}$$

where $\langle p | H_W^{PC} | \Sigma^+ \rangle = \alpha \bar{u}_p u_\Sigma$, $f_\pi = 133$ MeV.

Sum rules for α follow from the correlator (1) generalized to the case of the external field with the momentum k, $k^2 = 0$. In the first order in G_F we have

$$K(q_1, q_2) = \int \exp(iq_2 y - iq_1 x) \langle 0 | T\{\eta_p(y) \\ H_W(0) \bar{\eta}_\Sigma(x)\} | 0 \rangle dx\, dy \tag{8}$$

Let us introduce convenient notations

(8)

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$$K(q_1, q_2) = G_F c s \sqrt{2} \frac{\pi^4}{(2\pi)^{10}} \sum_{i=0}^{3} k_i(s_1, s_2) T_i \tag{9}$$

$$c = \cos\theta_c, \quad s = \sin\theta_c$$

where T_i are the four independent γ-matrix structures

$$T_0 = \not{q}_1 - \not{q}_2, \quad T_1 = \not{q}_1 + \not{q}_2, \quad T_2 = \frac{1}{2}(\not{q}_1 \not{q}_2 - \not{q}_2 \not{q}_1), \\ T_3 = -1 \tag{10}$$

Representing k_i in the form of the double dispersion relation

$$k_i(s_1, s_2) = \int_0^\infty \int_0^\infty \frac{\rho_i(u_1, u_2)\, du_1 du_2}{(u_1 + s_1)(u_2 + s_2)} \tag{11}$$

we can rewrite the double Laplace transformation of k_i in the form:

$$k_i(t_1, t_2) \equiv L_1 L_2 k_i(s_1, s_2) = \int_0^\infty \int_0^\infty \rho_i(u_1, u_2) \cdot \\ \cdot \exp\left(-\frac{u_1}{t_1} - \frac{u_2}{t_2}\right) \frac{du_1}{t_1} \frac{du_2}{t_2} \tag{12}$$

Phenomenological expressions for k_i follow from $K(q_1, q_2)$ saturated by the real hadronic states:

$$K(q_1, q_2) = \tilde{\beta}_p \tilde{\beta}_\Sigma (-\gamma_5) \frac{i}{\not{q}_2 - m_p} \alpha \frac{i}{\not{q}_1 - m_\Sigma} \gamma_5 + \\ + \text{higher states} \tag{13}$$

where $(2\pi)^2\langle 0|\eta_B|B\rangle = \tilde{\beta}_B\gamma_5 u_B$, $B = p,\Sigma$ and baryon bispinors are normalized as $\bar{u}u = 2m$. Applying the transformation L_1L_2 to $k_i(S_1,S_2)$ we suppress the higher states contribution in (13). Neglecting this contribution for the first time we have for the amplitude under consideration

$$A(\Sigma_0^+) = \frac{cs}{64\pi^2 m_\pi^2 f_\pi \tilde{\beta}_p\tilde{\beta}_\Sigma} t_1t_2\exp\left(\frac{m_\Sigma^2}{t_1}+\frac{m_p^2}{t_2}\right).$$

$$\cdot\begin{cases}\dfrac{2}{m_\Sigma+m_p}k_1(t_1,t_2), \text{ the sum rule I} \\ k_2(t_1,t_2), \text{ the sum rule II} \\ \dfrac{2}{(m_\Sigma+m_p)^2}k_3(t_1,t_2), \text{ the sum rule III}\end{cases} \tag{14}$$

The sum rule for the structure $T_0 = \mathcal{A}_1 - \mathcal{A}_2$ turns out to be useless for calculations; really, it is easy to check that k_0 is saturated mainly by the transitions $\frac{1}{2}^- \to \frac{1}{2}^+$, $\frac{1}{2}^+ \to \frac{1}{2}^-$, not by $\frac{1}{2}^+ \to \frac{1}{2}^+$.

The structure of the nonleptonic Hamiltonian H_W depends on the typical quark momentum Q_q. After substitution of H_W into (8) and Laplace transformation Q_q^2 is substituted by M^2, some scale of the order of t_1, t_2. Since $M \sim 1GeV \sim m_c$ we can use the result of refs. [13,14]

$$\begin{aligned} H_W(M) &= \sqrt{2}G_F cs(C_-(M)\mathcal{O}_- + C_+(M)\mathcal{O}_+) \\ \mathcal{O}_- &= (\bar{d}_L\gamma_\mu u_L)(\bar{u}_L\gamma^\mu s_L) - (\bar{u}_L\gamma_\mu u_L)(\bar{d}_L\gamma^\mu s_L) \\ \mathcal{O}_+ &= (\bar{d}_L\gamma_\mu u_L)(\bar{u}_L\gamma^\mu s_L) + (\bar{u}_L\gamma_\mu u_L)(\bar{d}_L\gamma^\mu s_L) \\ C_\pm(M) &= (\alpha_s(M)/\alpha_s(m_W))^{\lambda_\pm} \end{aligned} \tag{15}$$

where $\lambda_- = -2\lambda_+ = 4/9$ for the four quark flavours.

Our choice of the baryon currents corresponds to

$$\eta_p = (u^a C\gamma_\mu u^b)\gamma^\mu d^c\varepsilon_{abc}, \quad C = \gamma_0\gamma_2 \tag{16}$$

while the other currents are connected with η_p by $SU(3)$ transformations. The sum rules for the baryon residues $\tilde{\beta}_B$ into these currents (with exception of the case of Λ) were obtained in refs. [3-5]. They are presented in Appendix A.

Throughout the paper Ψ denotes the quark field with the flavour ($\alpha,\beta,\gamma,\ldots$ = 1,2,3) running over the three light flavours u, d, s, colour (a,b,c, ... = 1,2,3) and bispinor (i,k,l, ... = 1,2,3,4) indices, q means a definite light flavour (u_i^a, d_i^a or s_i^a). Expressions like $\bar{\Psi}\Psi$ imply the summation over all the indices while those like $\Psi\bar{\Psi}$ do not, i.e.

$$\bar{\Psi}\Psi \equiv \bar{\Psi}^a_{\alpha i}\Psi^a_{\alpha i}, \quad \Psi\bar{\Psi} \equiv \Psi^a_{\alpha i}\bar{\Psi}^b_{\beta k} \tag{17}$$

3. The weak interaction at short distances

Here we consider the graphs containing the weak vertex at short distances.

When calculating $k_i(t_1,t_2)$ in $SU(3)$ limit we take into account the following VEV's:

$$\langle I\rangle = 1,\ \langle\Psi\bar{\Psi}\rangle,\ \langle G^2\rangle,\ \langle\Psi G\bar{\Psi}\rangle,$$
$$\langle\Psi\bar{\Psi}\Psi\bar{\Psi}\rangle,\ \langle\Psi\bar{\Psi}\Psi\bar{\Psi}G\rangle \tag{18}$$

where the two last VEV's are estimated with the help of the factorization hypothesis, i.e. by vacuum insertion. Account of the anomalous dimensions of the operators is made as in refs. [3,4]; the anomalous dimension of an operator averaged by factorization is put to be equal to the sum of the anomalous dimensions of factors. Calculation of diagrams including emission of the soft gluons can be most conveniently performed in the fixed point gauge [6]. We place the coordinate origin either in H_W or, to check the answer, in $\bar{\eta}_\Sigma$ or η_p. Considering first the antisymmetric part of H_W (operator O_-) we get

$$k_i \equiv C_-(M)\tilde{k}_i,$$

$$
\tilde{k}_1 = \frac{6}{5}\frac{t_1^2 t_2^2}{t_1+t_2} L_M^{-4/9} + \frac{1}{6} b \frac{(t_1-t_2)^2}{t_1+t_2} L_M^{-4/9} + \frac{16}{3} a^2 L_M^{+4/9} \cdot \left[\frac{t_1^2+t_2^2}{2t_1t_2} + \frac{t_1t_2}{(t_1+t_2)^2}\right] + \frac{1}{9} a_g a \left[\frac{8}{t_1+t_2} - \frac{13}{2}\frac{t_1+t_2}{t_1t_2} - 13\frac{t_1^3+t_2^3}{t_1^2t_2^2} - 5I(0)\right], \tag{19}
$$

$$
\tilde{k}_2 = \frac{8}{3} a \frac{t_1t_2}{t_1+t_2} + a_g L_M^{-4/9}\left[\frac{1}{3} - \left(\frac{t_1-t_2}{t_1+t_2}\right)^2\right],
$$

$$
\tilde{k}_3 = 4a \frac{t_1^2+t_2^2}{(t_1+t_2)^2} t_1 t_2 + a_g \cdot 0
$$

where $I(Q^2)$ at $-k^2 = Q^2 \gtrsim 1\,GeV^2$ takes the form

$$
I(Q^2) = \int_0^\infty \frac{dt_3}{t_3} \frac{(t_1+t_2)\exp(-Q^2/t_3)}{(t_1+t_2+t_3)^2} \tag{20}
$$

and the following notations are introduced

$$
\begin{aligned}
L_M &= \frac{\alpha_s(\mu_0)}{\alpha_s(M)} = \frac{\ln M/\Lambda}{\ln \mu_0/\Lambda} \\
a &= -(2\pi)^2\langle\bar{q}q\rangle = 0.546\,GeV^3 \\
a m_0^2 &\equiv a_g = \langle i g_s \bar{q}\sigma_{\mu\nu} t^n q\, G^n_{\mu\nu}\rangle \\
b &= \langle g_s^2 G^n_{\mu\nu} G^n_{\mu\nu}\rangle \\
\sigma_{\mu\nu} &= \frac{1}{2}(\gamma_\mu\gamma_\nu - \gamma_\nu\gamma_\mu)
\end{aligned} \tag{21}
$$

$2t^n$ are the Gell-Mann matrices acting in colour space, and all the VEV's are normalized at a momentum μ_0.

Expression (20) for $I(Q^2)$ is valid at $Q^2 \gtrsim 1\ GeV^2$ and diverges logarithmically at $Q^2 \to 0$. However, at small Q^2 configurations in which the weak interaction occurs at large distances can be important in calculation of the considered quark-gluonic correction. Following the general idea of the operator expansion we must systematically separate the region

of large distances and include it in the definition of matrix elements of some (nonlocal) operators; small distances are treated perturbatively. Here we shall give an estimate of the large distance contribution into the considered correction using the model of the lowest state dominance in the four-quark channel in correlator at $Q^2 = 0$.

To find the residue of the lowest state pole in $I(Q^2)$ let us apply the Laplace transformation in Q^2 to $I(Q^2)$ with a parameter t_3 and then put for the rough estimate $t_3 = m_4^2$ where m_4 is a mass of the lowest four-quark state. Then we approximate $I(Q^2)$ by this pole at $Q^2 = 0$. The result reads

$$
I(0) = \frac{t_1+t_2}{(t_1+t_2+t_3)^2}\frac{t_3}{m_4^2}\exp\frac{m_4^2}{t_3} = \frac{t_1+t_2}{(t_1+t_2+m_4^2)^2}e^1 = \frac{e^1}{t_1+t_2} + O\left(\frac{1}{(t_1+t_2)^2}\right) \tag{22}
$$

Corresponding contribution into the correction due to $\langle\psi\bar{\psi}\psi G\bar{\psi}\rangle$ does not exceed 10%.

Of the VEV's considered (see (18)) only $\langle\psi\bar{\psi}\psi G\bar{\psi}\rangle$ gives contribution into the matrix elements of the symmetrical operator O_+ (in the graphs like that depicted in fig. 3). When calculating this contribution we reveal an uncertainty in sign connected with arising a difference of the terms of the same order of magnitude determined mainly by the large distance dynamics at $Q^2 \to 0$. However, we can say that this contribution is of the order of a few percents of the whole amplitude. A large value for the considered matrix element of O_+ would mean a violation of the Pati-Woo theorem [10], which holds experimentally with an accuracy $\simeq$5%. Therefore we neglect O_+ in this section.

Higher states contribution is accounted for in (13) in a model way. It is suggested that transferring the sum over the higher states from LHS to RHS of the sum rules reduces to limitation of the integration in the double dispersion relation (11) in the RHS by some region Ω. It seems the most natural to introduce a duality interval S_0 in each of the baryonic channels, i.e. to put Ω to be the square:

$$\Omega = \{(u_1,u_2) \mid 0 < u_i < s_0 \,,\; i = 1,2\} \tag{23}$$

Accepting this model let us give an estimate of $A(\Sigma_0^+)$ putting $t_1 \,,\, t_2 \to \infty$. Equally with an asymptotic loop the diagrams with iterations of the quark condensate $\langle \psi\bar{\psi} \rangle$ usually give a considerable contribution due to the loss of the small loop geometrical factors $(16\pi^2)^{-1}$. In our case the main term in k_1 is that proportional to $\langle \psi\bar{\psi}\psi\bar{\psi} \rangle$ and it is that proportional to $\langle \psi\bar{\psi} \rangle$ in k_2, k_3. As a result, we have in SU(3) limit from (14)

$$A(\Sigma_0^+) = \frac{c s C_-(\sqrt{s_0})}{64\pi^2 m_\pi^2 f_\pi} \begin{cases} \frac{16}{3} L_{\sqrt{s_0}}^{+4/9} \frac{a^2 s_0^2}{m \tilde{\beta}^2} \\ \frac{8}{9} \frac{a s_0^3}{\tilde{\beta}^2} \\ \frac{1}{6m^2} \frac{a s_0^4}{\tilde{\beta}^2} \end{cases} \tag{24}$$

At the same time the sum rules for the baryon residues give

$$\tilde{\beta}^2 = \begin{cases} L_{\sqrt{s_0}}^{-4/9} \frac{1}{12} s_0^3 \\ \frac{1}{2} a s_0^2 \end{cases} \tag{25}$$

Combining (24) and (25) so as to cancel the s_0 independence in (24) as far as it is possible we get

$$A(\Sigma_0^+) = \frac{cs}{3} \frac{f_\pi}{m_u + m_d} C_-(\mu_0) \begin{cases} 1 \\ 1 \\ \frac{3}{16} \frac{s_0}{m^2} \end{cases} \tag{26}$$

Here the known result of PCAC is used [12]

$$\langle \bar{q}q \rangle = -\frac{1}{2} \frac{f_\pi^2 m_\pi^2}{m_u + m_d} \tag{27}$$

Thus, we see that the sum rules I and II lead in a consistent way to the estimate

$$A(\Sigma_0^+) = \frac{cs}{3} \frac{f_\pi}{m_u + m_d} C_-(\mu_0) \simeq 2 \tag{28}$$

at $m_u + m_d$ = 11 MeV, cs = 0.215, Λ = 0.1 GeV, $\alpha_s(\mu_0)$ = = 0.7. As for the sum rule III it possesses the largest sensitivity to the continuum contribution due to the high power of t_1, t_2 in the main term. Here we can state a consistency with I and II only up to a factor of 2 (at the typical values $s_0 \simeq$ 2÷3 GeV^2).

Hereafter we consider only the sum rule I which possesses a comparatively small sensitivity to the continuum (as it is seen from (26)). Besides, the structure $T_1 = q_1 + q_2$ survives at $q_1 = q_2$ which is important for the analysis of the weak interaction at large distances.

Finally, let us present the corrections to k_1 taking into account the SU(3) violating operators

$$m_s \psi\bar{\psi},\; m_s \psi G \bar{\psi},\; u\bar{u} - s\bar{s},\; (u\bar{u} - s\bar{s}) G \tag{29}$$

While in SU(3) limit we have had

$$A(\Sigma_0^+) = -A(\Lambda_-^0)\sqrt{3} = \frac{\sqrt{3}}{2} A(\Xi_-^-) \sim k_1 \tag{30}$$

now three functions k_{1B}, $B = \Sigma, \Lambda, \Xi$ arise for the decays Σ, Λ, Ξ respectively,

$$\begin{aligned} & k_{1B} = C_-(M) \tilde{k}_{1B}\,, \\ & \tilde{k}_{1\Sigma} - \tilde{k}_1 = \frac{2}{3} m_s a L_M^{-4/9} \left[3 \frac{t_1^2 t_2}{(t_1+t_2)^2} - 2 \frac{t_2^2}{t_1+t_2} \right] - \\ & \quad - \frac{1}{4} m_s a_g L_M^{-8/9} \left[1 + \frac{t_1}{t_2} - 4 \frac{t_1 t_2}{(t_1+t_2)^2} \right] + \\ & \quad + \frac{8}{3} a^2 f L_M^{+4/9} \frac{t_1 t_2}{(t_1+t_2)^2} + \frac{1}{18} a_g a f \left(\frac{20 - 3e'}{t_1+t_2} + \right. \end{aligned} \tag{31}$$

$$+ \frac{5}{2}\frac{t_1+t_2}{t_1t_2}\Big),$$

$$\tilde{k}_{1\Lambda}-\tilde{k}_{1\Sigma} = 8m_s a L_M^{-4/9} t_2 - \frac{1}{3} m_s a_g L_M^{-8/9}\Big(\frac{3t_1^2+11t_1t_2+14t_2^2}{t_1t_2} - 6\frac{t_2}{t_1+t_2}\Big) + \frac{16}{3}a^2 f L_M^{+4/9}\frac{t_2}{t_1} -$$

$$- \frac{2}{9}a_g a f\Big(\frac{-t_1^2+9t_1t_2+13t_2^2}{t_1^2t_2} + \frac{12+2e^1}{t_1+t_2}\Big),$$

$$\tilde{k}_{1\Xi}-\tilde{k}_{1\Sigma} = -\frac{4}{3} m_s a L_M^{-4/9}\Big(t_1-2t_2+4\frac{t_1t_2}{t_1+t_2}\Big) +$$

$$+ \frac{1}{6} m_s a_g L_M^{-4/9}\Big(\frac{8t_1^2+13t_1t_2-3t_2^2}{t_1t_2} + 6\frac{t_2}{t_1+t_2}\Big) +$$

$$+ \frac{8}{3}a^2 f L_M^{+4/9}\frac{t_1^2+2t_2^2}{t_1t_2} - \frac{1}{9}a_g a f\Big(\frac{12+2e^1}{t_1+t_2} +$$

$$+ \frac{9}{t_2} + \frac{19}{t_1} + 13\frac{t_1}{t_2^2} + 26\frac{t_2}{t_1^2}\Big)$$

where $f = \langle(\bar{s}s-\bar{u}u)\rangle/\langle\bar{u}u\rangle \approx f_g =$
$= \langle(\bar{s}Gs-\bar{u}Gu)\rangle/\langle\bar{u}Gu\rangle.$

4. Interaction of a variable external field with the soft quarks

Here we consider a possibility to express a result of the double Laplace transformation L_1L_2 at $t_1 = t_2$ through that of the ordinary one L_t in $S=S_1=S_2$ for the diagram relevant to the interaction of an external field with the soft quarks (i.e. at large distances). Namely, we give some arguments in favour of (5) paying a special attention to the case of the weak field.

Firstly, this model looks quite reasonable from the viewpoint of intuitive approach: once interaction is soft, it cannot appreciably change the invariant mass squared S_1 of incoming three-quark state. This circumstance is just reflected

by the δ-function in (5). Secondly, the double spectral density of this kind can be obtained by explicit calculation of the graphs of figs. 5a, b* containing the weak vertex at short distances. This fact can be viewed as an indication to that the same property holds for the diagrams of a similar structure although relevant to the weak interaction at large distances (see fig. 7). Thirdly, our hypothesis is supported by an analogy with the polarization operators in external fields coupled to the quark bilinears. Here we discuss the corresponding sum rules in more detail.

Let us first consider the case of a constant external field to which we want to reduce the case of the variable one. Now single pole terms not cancelled by L_t are to be taken into account in the LHS of the sum rules. The latter take the general form [9,11]

$$\tilde{\beta}^2\Big(\frac{A}{t^2}+\frac{B}{t}\Big)\exp\Big(-\frac{m^2}{t}\Big) + \sum_i \tilde{\beta}_i^2\Big(\frac{A_i}{t^2}+\frac{B_i}{t}\Big)\cdot \exp\Big(-\frac{m_i^2}{t}\Big) = \sum_n C_n t^{n-2} \qquad (32)$$

Here A is a baryon-baryon-external field coupling sought for, B parametrizes the single pole terms, i labels the excited states. The model of continuum suggested in refs. [9, 11] allows one to rewrite (32) in the form

$$\tilde{\beta}^2(A+Bt) = \exp\Big(\frac{m^2}{t}\Big)\sum_n C_n t^n E_n\Big(\frac{S_0}{t}\Big) \qquad (33)$$

where

---

*) Contribution of the graphs like that depicted in fig. 5c relevant to the perturbative non-diagonal quark mass insertion turns out to be several times smaller than that of the graphs of figs. 5a,b. In turn, the latter is only a few percents radiative correction to the whole amplitude. Therefore the graphs considered are omitted in the subsequent analysis.

$$E_n(x)=\begin{cases}\frac{1}{(n-1)!}\int\limits_0^x y^{n-1}e^{-y}dy, & n\geq 1\\ 1, & n\leq 0\end{cases} \tag{34}$$

It was shown in ref. [11] that in the case of vector field C_n reproduce the coefficients in the expression for $\tilde{\beta}^2$. As a result, the ansatz (33) presents the only possibility to ensure the conservation of vector current, $A=g_V=1$, S_0 being the continuum threshold in the sum rules for $\tilde{\beta}^2$. In general, A and B are determined by an asymptotic behaviour of the RHS of (33) and so they crucially depend on the model of continuum used. The results of refs. [9,11] show that the ansatz (33) presents quite reasonable choice of this model for it leads to the consistency with the experiment.

Further, a dependence of the RHS on S_0 and t imposes some limitations on the form of the double spectral density $\rho(S_1,S_2)$. Really, we have for the contribution of an arbitrary term in (33) to this density

$$C_n t^{n-2}E_n\left(\frac{S_0}{t}\right)=L_t\int\limits_{\Omega}\frac{\rho(u_1,u_2)\,du_1du_2}{(u_1+s_1)(u_2+s_2)} \tag{35}$$

The integration region Ω is parametrized by S_0. In principle, one must write out in the RHS of (35) also single pole subtraction terms not cancelled by L_t. However, the only necessary subtraction in the diagram of interest of fig. 4 is to be done in the subdiagram I; the shaded blob does not add any new divergences for it survives, by definition, only at small momenta p_1^2, p_2^2 and in the case $p_1=p_2=k=0$ it reduces to the finite matrix element of a definite (nonlocal) operator. Corresponding subtraction term is polinomial in S_1, S_2 and thus it is cancelled by L_t.

It is shown in Appendix B that (35) leads to

$$\rho(s_1,s_2)=C_n\frac{s_1^{n-1}}{(n-1)!}\delta(s_1-s_2) \tag{36}$$

If the spectral density of a function $f(S_1,S_2)$ is proportional to $\delta(S_1-S_2)$, one has for the Laplace transformations of $f(S_1,S_2)$:

$$L_1L_2f(s_1,s_2)=\frac{t^2}{t_1t_2}L_tf(s,s)$$
$$L_1L_2f(s_1,s_2)=\frac{1}{4}L_tf(s,s)\quad \text{at}\quad t_1=t_2=2t \tag{37}$$

5. The weak interaction at large distances

Here we consider the graphs determined by the dependence of the VEV's on G_F.

Namely, we have taken into account the following VEV's different from zero due to the weak interaction and contributing to the γ-matrix structure $\not{q}$ of interest:

$$\langle\nabla_\mu d\bar{s}\rangle,\ \langle\nabla_\mu d\bar{s}G\rangle\ (\langle\nabla_\mu\nabla_\nu\nabla_\lambda d\bar{s}\rangle),\ \langle d\bar{s}q\bar{q}\rangle \tag{38}$$

where the last VEV is studied in the framework of factorization

$$\langle d\bar{s}q\bar{q}\rangle\approx\langle d\bar{s}\rangle\langle q\bar{q}\rangle \tag{39}$$

These VEV's generate the graphs presented in figs. 7a,b,c,d in which the two quarks $\bar{s}$ and d weakly interact at large distances. As for the graphs in which the four (valence) quarks $\bar{s}$, d, $\bar{u}$, u interact at large distances, their contribution was estimated in sect. 3 (see the consideration after (19)-(21)) and turned out to be negligible for the given orders of the power expansion.

The VEV's (38) can be differentiated in G_F with the help of the equations of motion

$$i\not{\nabla}q=m_q q+\frac{\partial H_w}{\partial\bar{q}} \tag{40}$$

and thus reduced to the pure QCD VEV's of some local operators. For example,

$$<d_i^a \bar{s}_k^b> = -\frac{1}{12} <0| \frac{\bar{s}\frac{\partial H_w}{\partial \bar{d}} - \frac{\partial H_w}{\partial s} d}{m_s - m_d} |0> \delta^{ab} \delta_{ik} \qquad (41)$$

$$<\nabla_\mu d_i^a \bar{s}_k^b> = \frac{i}{48} <\bar{s} \frac{\partial H_w}{\partial \bar{d}}> (\gamma_\mu)_{ik} \delta^{ab} \qquad (42)$$

These VEV's are determined at the typical momentum $M \simeq m_c$. Let us express them through the VEV's of the operators determined at a momentum μ sufficiently small to allow us to calculate these new VEV's by the vacuum insertion. Some new four-fermion operators appear at such the renormalization which include both the left- and right-handed quark fields (contrary to the original operators normalized at m_c). Just these operators give a non-zero result at the vacuum averaging. In the first order in α_s they arise due to the annihilation graphs of figs. 6a,b.

Let us consider the four-fermion operator entering (41) in more detail. It possesses the following structure

$$(\bar{d}_L \Gamma d_L - \bar{s}_L \Gamma s_L)(\bar{u}_L \Gamma u_L) \qquad (43)$$

where $\Gamma \otimes \Gamma = \gamma_\mu \otimes \gamma^\mu$ or $\gamma_\mu t^n \otimes \gamma^\mu t^n$. Its renormalization due to the graphs of fig. 6a leads to the operator

$$(\bar{d}_L \gamma_\mu t^n d_L - \bar{s}_L \gamma_\mu t^n s_L)(\bar{\psi}_R \gamma^\mu t^n \psi_R) \ln \frac{m_c^2}{\mu^2} \qquad (44)$$

which being sandwiched between the vacuum states is proportional to

$$(<\bar{d}d>^2 - <\bar{s}s>^2) \ln \frac{m_c^2}{\mu^2} \qquad (45)$$

The graphs of fig. 6b add the following term to (45)

$$<\bar{u}u>^2 I\left(\frac{m_s^2}{\mu^2}\right), \quad I(\lambda) = 6\int_0^1 d\alpha\, (\alpha - \alpha^2) \cdot$$

$$\cdot \ln\left(1 + \frac{\lambda}{\alpha - \alpha^2}\right) = 6\lambda + O(\lambda^2) \quad (\text{at } m_d = 0) \qquad (46)$$

This term does not contain formally large logarithm while (45) does. Besides, (46) vanishes in the chiral limit $m_{u,d,s} \to 0$. However, these factors diminishing the ratio of (46) to (45) are compensated, firstly, by the smallness of the value $(<\bar{d}d>^2 - <\bar{s}s>^2)$ vanishing in SU(3) limit, secondly, by the large numerical factor 6 in (46). If we start from $\mu \gg m_s$ then the ratio of (46) to (45)

$$\frac{3}{2} \, \frac{m_s^2}{-f} \, \frac{1}{\mu^2 \ln \frac{m_c}{\mu}} \qquad (47)$$

reaches unity at $\mu \simeq$ 0.5 GeV. Here $f = <(\bar{s}s - \bar{u}u)>/ /<\bar{u}u> \simeq -(0.15 \div 0.20)$. Thus, not only u-, c-, but also s-, d-quarks being closed in loop can give a significant contribution to the VEV's considered.

Let us present expressions for the VEV's of interest

$$<d_i^a \bar{s}_k^b> = \frac{G_F c s \sqrt{2}}{27} <\bar{q}q>^2 \cdot 4 C_{1R} \frac{-f}{m_s} \delta_{ik} \delta^{ab}$$

$$<\nabla_\mu d_i^a \bar{s}_k^b> = \frac{G_F c s \sqrt{2}}{27} <\bar{q}q>^2 C_{2R} (i\gamma_\mu)_{ik} \delta^{ab}$$

$$<\nabla_\mu d_i^a \bar{s}_k^b g_s G^n_{\rho\lambda}> = \frac{G_F c s \sqrt{2}}{24 \cdot 16} <\bar{q}q>^2 m_0^2 C_{3R} \cdot$$

$$\cdot \frac{1}{2} (\gamma_\rho \gamma_\mu \gamma_\lambda - \gamma_\lambda \gamma_\mu \gamma_\rho)_{ik} (t^n)^{ab} \qquad (48)$$

$$x^\mu x^\nu x^\lambda <\nabla_\mu \nabla_\nu \nabla_\lambda d_i^a \bar{s}_k^b> = \frac{G_F c s \sqrt{2}}{24 \cdot 16} <\bar{q}q>^2 m_0^2 \cdot$$

$$\cdot C_{3R} \cdot 2 x^2 \cdot (i \not{x})_{ik} \delta^{ab}$$

where with taking into account the nonlogarithmic terms

$$C_{1R} = \frac{C_- + C_+}{12\pi}\,\alpha_s(\mu)\left[\ln\frac{m_c}{\mu} + \frac{5}{6} + \frac{3}{-2f}\, I\left(\frac{m_s^2}{\mu^2}\right)\right],$$

$$C_{2R} = \frac{C_- + C_+}{12\pi}\,\alpha_s(\mu)\left[\ln\frac{m_c}{\mu} + \frac{5}{6} - \frac{3}{2}\, I\left(\frac{m_s^2}{\mu^2}\right)\right],$$

$$C_{3R} = \frac{2C_- + C_+}{18\pi}\,\alpha_s(\mu)\left(\ln\frac{m_c}{\mu} + \frac{5}{6}\right) - \qquad\qquad$$

$$- \frac{11C_- + 5C_+}{36\pi}\,\alpha_s(\mu)\, I\left(\frac{m_s^2}{\mu^2}\right) \qquad (49)$$

$$C_- \equiv C_-(m_c), \quad C_+ \equiv C_+(m_c)$$

At calculations we have also summed up using the renormalization group (RG) technique the terms of the kind $\alpha_s^n \ln^n m_c/\mu$ appearing in C_{1R}, C_{2R} in the next orders of perturbation theory. Operators (43) entering (41) and contributing to C_{1R} are connected by means of U -spin rotation with the following operators

$$(\bar{d}_L \Gamma s_L + \bar{s}_L \Gamma d_L)(\bar{u}_L \Gamma u_L) \qquad (50)$$

entering the effective $|\Delta S| = 1$, $|\Delta T| = 1/2$ weak Hamiltonian. So they have the same RG properties and

$$C_{1R}^{LLA} = -C_5(\mu) - \frac{3}{16} C_6(\mu) \qquad (51)$$

Expressions for $C_{5,6}(\mu)$ were obtained in ref. [15]. As for the operators entering (42) and contributing to C_{2R} only the flavour singlets are essential in SU(3) limit. They belong to the following set of operators closed under RG transformations

$$(\bar{\Psi}_L \Gamma \Psi_L)(\Psi_L \Gamma \Psi_L) + (L \to R)$$
$$(\bar{\Psi}_L \Gamma \Psi_L)(\bar{\Psi}_R \Gamma \Psi_R) \qquad (52)$$

Their RG properties are given in ref. [1]. At $\alpha_s(m_c)/\alpha_s(\mu) = 0.3$ taking into account the higher powers of logarithms provides the negative corrections of approximately 15% and 30% to the one-logarithmic terms in the expressions (49) for C_{1R}

and C_{2R}, respectively.

Calculation of the graphs of figs. 7a,b,c,d is performed at $q_1 = q_2 = q$ making use of the formula (37) for the double Laplace transformation in terms of the ordinary one. The final result at $t_1 = t_2 = 2t$ has the form

$$\left.\begin{array}{c} A(\Sigma_0^+) \\ -A(\Lambda_-^0)\sqrt{3} \\ A(\Xi_-^-)\sqrt{3}/2 \end{array}\right\} = \frac{2}{m_1 + m_2}\,\frac{1}{\tilde{\beta}_1\tilde{\beta}_2}\,\frac{cs}{16\pi^2 m_\pi^2 f_\pi}\,\cdot$$

$$\cdot \exp\left(\frac{m_1^2 + m_2^2}{2t}\right)\cdot t^2 \left\{\begin{array}{l} k_{1\Sigma} + k_{5\Sigma} \\ k_{1\Lambda} + k_{5\Lambda} \\ k_{1\Xi} + k_{5\Xi} \end{array}\right. \qquad (53)$$

where k_{1B} are given in sect. 3 and

$$k_{5\Sigma} = \frac{32}{9} C_{2R}\, a^2 L_M^{-4/9} \eta^{-8/9} E_2(x) -$$
$$- \frac{2}{3} C_{3R}\, a^2 m_0^2 L_M^{-4/9} \eta^{-4/9} \frac{E_1(x)}{t}$$
$$k_{5\Lambda} = 3k_{5\Sigma} + \frac{256}{27} C_{1R}\, a^3 \frac{-f}{m_s} L_M^{+4/9} \eta^{-8/9} \frac{E_1(x)}{t} \qquad (54)$$
$$k_{5\Xi} = \frac{3}{4} k_{5\Sigma} + \frac{1}{4} k_{5\Lambda}$$
$$\eta = \frac{\alpha_s(\mu)}{\alpha_s(\mu_0)}, \quad x = \frac{s_0}{t}$$

The RG parameter η is needed to express the VEV's $\langle\bar{q}q\rangle$, $\langle\bar{q}Gq\rangle$, $\langle G^2\rangle$ and the mass m_s taken at μ through their values at μ_0 at which these values are considered (by definition of μ_0) as known ones. The usual choice of μ_0, μ satisfies

$$\alpha_s(\mu_0) = 0.7, \quad \alpha_s(\mu) = 1 \qquad (55)$$

The threshold s_0 in (54) is defined so that (s_0, s_0) is the point of intersection of the diagonal $s_1 = s_2$ with the boundary of Ω .

Varying s_0 we aim at the existence of a platou in the dependence of the amplitude of interest on t . Power corrections must not exceed, say, 30% of the main term in order that the whole calculation using a few first terms of the operator expansion might be reliable [1]. The dependences of $A(\Sigma_0^+)$ on t are depicted graphically in fig. 8 for three choices of s_0 . As it was mentioned in sect. 3, the term proportional to $\langle\psi\bar{\psi}\psi\bar{\psi}\rangle$ gives the leading contribution into the amplitude. The ratio of the next power correction due to $\langle\psi\bar{\psi}\psi G\bar{\psi}\rangle$ to this term is 30% at t = 1 GeV^2 and 20% at t = 1.5 GeV^2.

At the same time sum rules are the most sensitive to the lowest resonance contribution just at $t \simeq \overline{m^2} = (m_1^2+m_2^2)/2 \approx 1$ GeV. It is seen from fig. 8 that $A(t)$ is practically constant at $0.6 < t < 1.5$ GeV^2, s_0 = 2.15 GeV^2 and is equal to

$$A(\Sigma_0^+) = 1.44 \tag{56}$$

The typical scales of $\overline{m^2}$, s_0, t become approximately 1.5 times larger for the Ξ decays.

Using the triangle version for Ω

$$\Omega = \{(u_1,u_2) \mid u_1+u_2 < 2s_0,\ u_1>0,\ u_2>0\} \tag{57}$$

instead of the square one (23) enhances the theoretical values for $A(\Sigma_0^+)$, $A(\Lambda_-^0)$ slightly (by no more than 10%) but this enhancement can reach 25% for $A(\Xi_-^-)$. We stick to the square version for Ω since namely this model leads to the identity of the duality estimates (26) for the two different γ -matrix structures and, in addition, this model posseses a more transparent physical sense.

The results of the calculation are given in table 1 for

$$m_c = 1.25\,GeV,\ m_s = 0.15\,GeV,\ f = -0.2 \tag{58}$$

and for the three choices of m_0^2, Λ . The residues $\tilde{\beta}_B$ were calculated independently for each choice of the parameters making use of the sum rules given in Appendix A. The amplitudes A_1 obtained with the neglection of k_5 in the correlator

are also presented in the table. Index 5 serves to remind of the definitive role of the annihilation mechanism in k_5 leading to appearance of the operator structures analogous to O_5, O_6. It is seen that this mechanism contributes mainly to the Λ decays. This fact is connected with that the Born graph of fig. 7d yielding the dominant contribution to $k_{5\Lambda}$ (the last term in $k_{5\Lambda}$ (54)) does not contribute to $k_{5\Sigma}$.

6. Conclusion

Up to now the main difficulties usually encountered at calculations of the non-leptonic hyperon decays have been connected with a choice of the normalization point of the effective Hamiltonian and with a modelless accounting for different mechanisms on equal footing. We see that the QCD sum rules method proposes a way to surmount these difficulties to a considerable extent.

In conclusion, the author is grateful to V.L.Chernyak, I.B.Khriplovich, A.I.Vainshtein, O.V.Zhirov, A.R.Zhitnitsky and I.R.Zhitnitsky for numerous helpful discussions and reading the manuscript. Discussions with Prof. A.I.Vainshtein had led to understanding of the role of the long-distance dynamics.

Appendix A.

The sum rules for the baryon residues into the currents used have the form

$$2\tilde{\beta}_N^2 \exp\left(-\frac{m_N^2}{t}\right) = t^3 E_3(x) L_M^{-4/9} + \frac{b}{4} t E_1(x) L_M^{-4/9} + \frac{4}{3} a^2 \left(L_M^{+4/9} - \frac{m_0^2}{4t}\right) \equiv f(t) \qquad (1A)$$

$$2\tilde{\beta}_\Sigma^2 \exp\left(-\frac{m_\Sigma^2}{t}\right) = f(t) - 2 m_s a t E_1(x) L_M^{-4/9} \qquad (2A)$$

$$2\tilde{\beta}_\Lambda^2 \exp\left(-\frac{m_\Lambda^2}{t}\right) = f(t) + 2 m_s a \left(\frac{1}{3} - f\right) t E_1(x) L_M^{-4/9} + \frac{16}{9} a^2 L_M^{+4/9} f \qquad (3A)$$

$$2\tilde{\beta}_\Xi^2 \exp\left(-\frac{m_\Xi^2}{t}\right) = f(t) + \frac{4}{3} a^2 (2f + f^2) L_M^{+4/9} \qquad (4A)$$

$$x = \frac{S_0}{t}, \quad L_M = \ln\left(\frac{M}{\Lambda}\right) / \ln\left(\frac{\mu}{\Lambda}\right), \quad t = M^2$$

These are taken (with exception of (3A)) from refs. [3,4].

Appendix B.

Let us rewrite (35) in the form ($C_n = 1$)

$$t^{n-2} E_n\left(\frac{S_0}{t}\right) = \int_\Omega \frac{\rho(u_1,u_2)}{u_1 - u_2} \left[\exp\left(-\frac{u_2}{t}\right) - \exp\left(-\frac{u_1}{t}\right)\right] \frac{du_1 du_2}{t} \qquad (1B)$$

Everywhere the symmetrical part of $\rho(u_1,u_2)$ is of interest, i.e. we can put

$$\rho(u_1,u_2) = \rho(u_2,u_1) \qquad (2B)$$

without loss of generality. Let us divide Ω into two regions: a narrow band along the diagonal, $|u_1 - u_2| < \varepsilon$, and its complement. Then (1B) at $\varepsilon \to 0$ takes the form

$$\int_0^{S_0} \frac{u^{n-1}}{(n-1)!} \exp\left(-\frac{u}{t}\right) du = \int \left[2t \cdot V.p. \int_0^{\varphi(u_1,S_0)} \frac{\rho(u_1,u_2)}{u_2 - u_1} du_2 + \sum_{k=0}^{\infty} \frac{f_k(u_1)}{t^k}\right] \exp\left(-\frac{u_1}{t}\right) du_1 \qquad (3B)$$

where

$$(k+1)!\, f_k(u_1) = \int_{u_1-0}^{u_1+0} (u_1 - u_2)^k \rho(u_1,u_2)\, \chi_\Omega(u_1,u_2)\, du_2,$$

$$\chi_\Omega(u_1,u_2) = \begin{cases} 1, & (u_1,u_2) \in \Omega \\ 0, & (u_1,u_2) \notin \Omega \end{cases} \qquad (4B)$$

and equation $u_2 = \varphi(u_1, S_0)$ determines the boundary of Ω. An inversed Laplace transformation of (3B) results in

$$V.p. \int_0^{\varphi(u_1,S_0)} \frac{\rho(u_1,u_2)}{u_2 - u_1} du_2 = 0 \qquad (5B)$$

$$\left.\begin{aligned} f_0(u_1) &= \frac{u_1^{n-1}}{(n-1)!} \theta(S_0 - u_1)\theta(u_1) \\ f_k(u_1) &= 0, \quad k \geq 1 \end{aligned}\right\} \qquad (6B)$$

It follows from (5B) that $\rho(u_1,u_2) = 0$ at $u_1 \neq u_2$ hence

$$\rho(u_1,u_2) = \sum_{n=0}^{N} a_n \delta^{(n)}(u_1 - u_2) \qquad (7B)$$

Then accounting for (4B), (6B) we finally get

$$\rho(u_1,u_2) = \frac{u_1^{n-1}}{(n-1)!} \delta(u_1 - u_2) \qquad (8B)$$

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Figure captions

Fig. 1. Asymptotic loop. The dotted lines cut the diagram across the channels in which the lowest states are to be separated.

Fig. 2. A diagram relevant to the four-quark operator $\psi\bar{\psi}\psi\bar{\psi}$ in the operator expansion.

Fig. 3. A diagram giving rise to the operator $\psi\bar{\psi}\psi G\bar{\psi}$. The dotted circles denote two (ψ , $\bar{\psi}$) or three ($\psi, \bar{\psi}, G$) factorized vacuum fields.

Fig. 4. The graph relevant to an interaction of an external field of momentum k with the soft quarks.

Fig. 5. Graphs possessing the same topology as those calculated below (see fig. 7) but containing the weak vertex at short distances.

Fig. 6. Annihilation graphs.

Fig. 7. The diagrams relevant to a dependence of the VEV's studied on G_F . Symbol W means accounting for the weak interaction in the first order in G_F . Gross denotes the linear term in the expansion of vacuum fields in x_μ , the distance between the baryon currents.

Fig. 8. The amplitude $A(\Sigma_0^+)$ versus t for different thresholds S_0. Λ = 75 MeV, m_0^2 = 1 GeV^2.

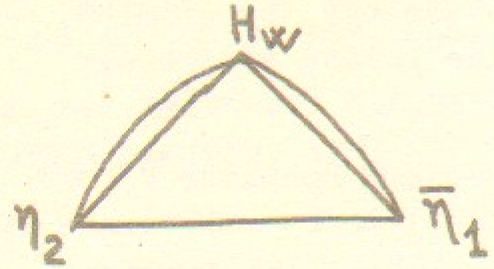

Fig. 1

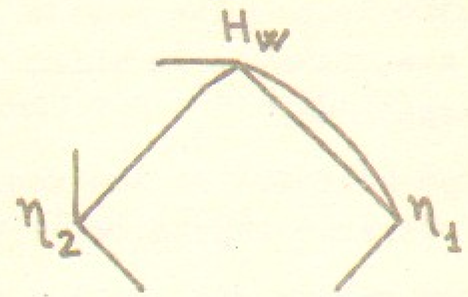

Fig. 2

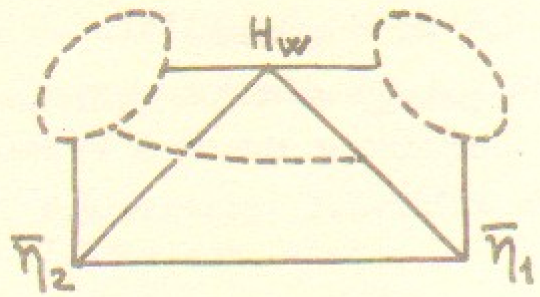

Fig. 3

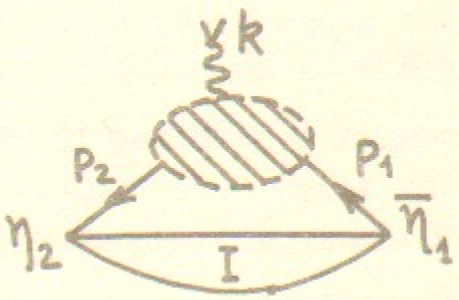

Fig. 4

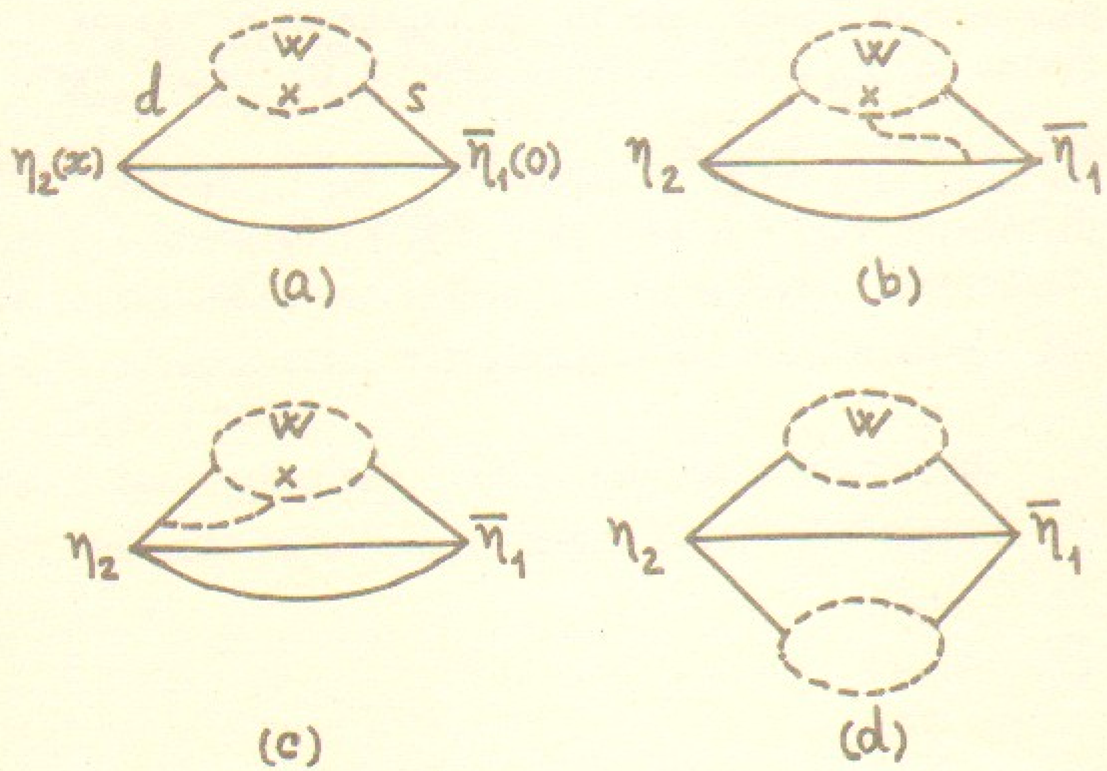

Fig. 7

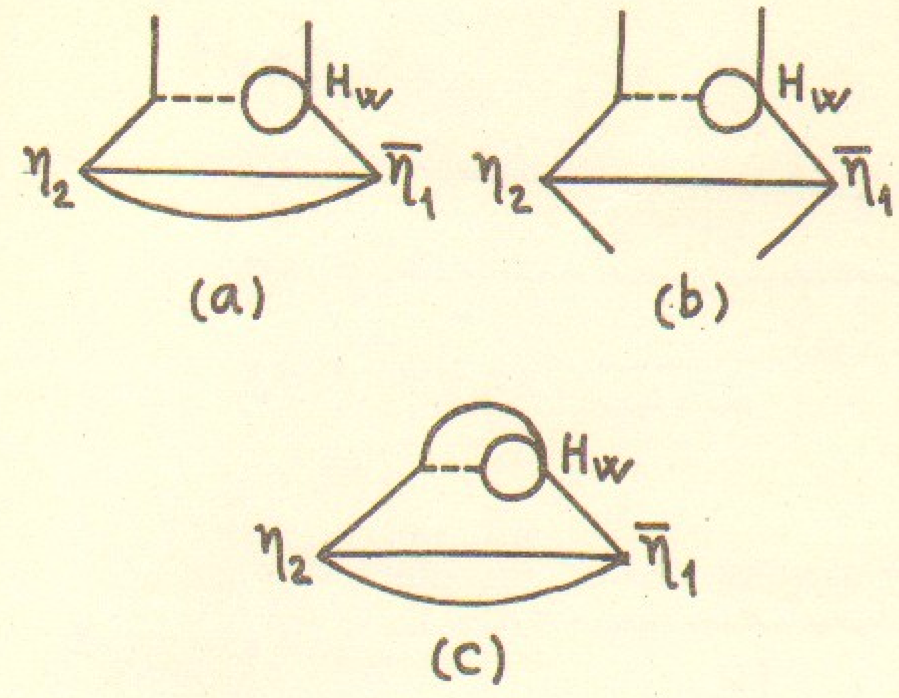

Fig. 5

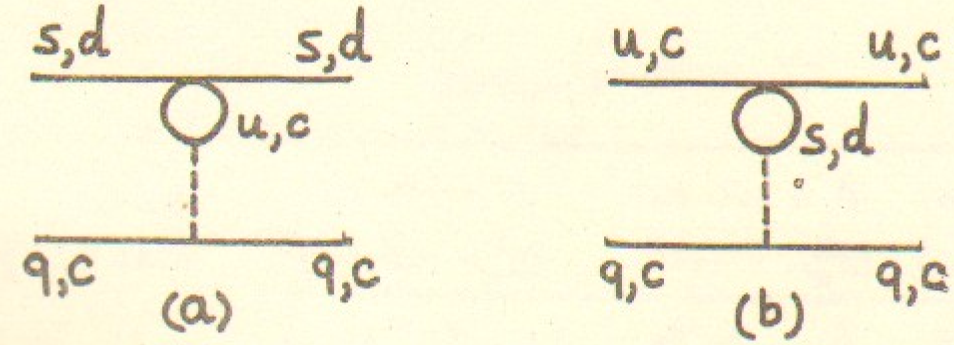

Fig. 6

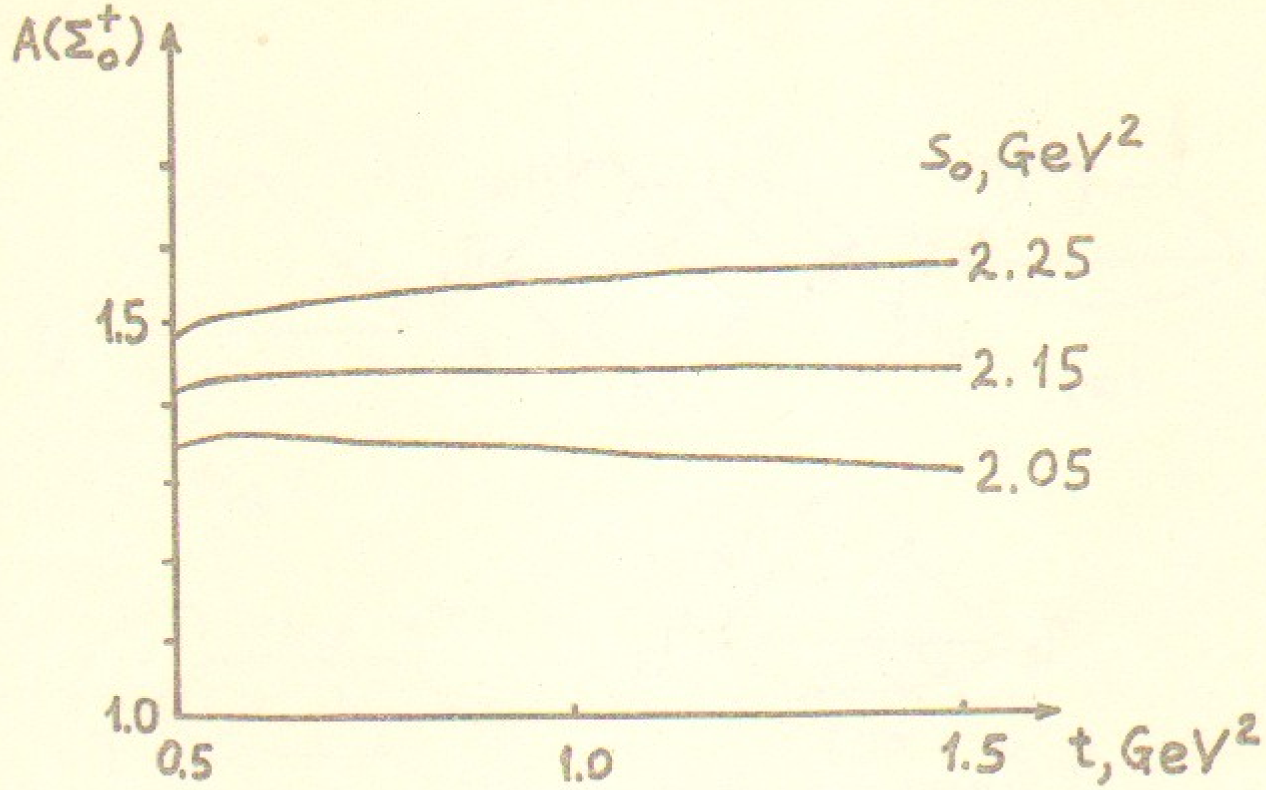

Fig. 8

Table1.

S-waves*

Mode	Λ = 75 MeV m_o^2= 1 GeV^2		Λ = 100 MeV m_o^2= 1 GeV^2		Λ = 100 MeV m_o^2= 0.8 GeV^2		Experiment [17]
	A_1	A_1+A_5	A_1	A_1+A_5	A_1	A_1+A_5	
Σ_o^+	1.44	1.44	1.35	1.39	1.65	1.71	1.48±0.05
Λ_-^o	0.73	1.35	0.66	1.22	0.87	1.37	1.48±0.01
Ξ_-^-	2.03	2.30	2.02	2.23	2.17	2.41	2.04±0.01

*) $A_1 + A_5$ means the full amplitude; amplitude A_1 is obtained neglecting the weak interaction at large distances.

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S - ВОЛНОВЫЕ НЕЛЕПТОННЫЕ РАСПАДЫ ГИПЕРОНОВ В КХД

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