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A.A. Ogloblin

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On the Σ^*- and Ξ^*-Hyperons Wave Functions

A.A. Ogloblin

Institute of Nuclear Physics
630090, Novosibirsk, USSR

ABSTRACT

The properties of the leading twist wave functions of the $\Sigma^*(1385)$ and $\Xi^*(1530)$-hyperons are investigated using the QCD sum rules. The model wave functions which satisfy the sum rules requirements are proposed. The asymptotic behaviour of Σ^*- and Ξ^*-hyperons electromagnetic formfactors and that of various «octet-decuplet» transition formfactors is found.

1. INTRODUCTION

In references [1, 2] the method was developed which allows one to calculate the moments of nonperturbative hadronic wave functions by using QCD sum rules [3]. Below using this method we continue to investigate the leading twist wave function properties of baryons entering the $\Delta(1232)$ decuplet. The leading twist wave functions properties of the $\Delta(1232)$- and $\Omega^-(1660)$-hyperons were investigated in [4], and the following model wave functions were proposed:

$$\varphi_\Delta(x)=\varphi_{ac}(x)\,[4.2\,(x_1^2+x_3^2)+2.52x_2^2-6.72x_1x_3+0.42x_2(x_1+x_3)]\,, \qquad (1)$$

$$\varphi_{\Omega^-}(x)=\varphi_{ac}(x)\equiv 5!\,x_1x_2x_3\,, \qquad (2)$$

The values of these wave functions at the origin were found there also:

$$f_\Delta=1.2\cdot10^{-2}\,\mathrm{GeV}^2, \quad f_\Omega=1.6\cdot10^{-2}\,\mathrm{GeV}^2. \qquad (3)$$

The wave function (WF) $\Phi(x_i)$ is the fundamental object of the theory and it describes the distribution of quarks in the hadron in longitudinal momentum fractions $0\leqslant x_i\leqslant 1$, $\Sigma x_i=1$ (at $p_z\to\infty$).

The Δ^+-hyperon leading twist WF is determined by the matrix element of the three-local operator, which is analogous to (5) with change $s\to d$; has the same properties of symmetry (7) and due to the SU(2)-symmetry $T_\Delta(x_1,\,x_2,\,x_3)=\varphi_\Delta(x_1,\,x_3,\,x_2)$. Δ^+-state can be written in the form:

$$|\Delta^+(\lambda=1/2)\rangle = \frac{f_\Delta}{\sqrt{3}}\int_0^1 \frac{d_3x}{4\sqrt{24x_1x_2x_3}}\,\varphi_\Delta(x)\times$$

$$\times\{|u^\uparrow(x_1)\,u^\downarrow(x_2)\,d^\uparrow(x_3)\rangle + |u^\uparrow(x_1)\,d^\downarrow(x_2)\,u^\uparrow(x_3)\rangle +$$
$$+ |d^\uparrow(x_1)\,u^\downarrow(x_2)\,u^\uparrow(x_3)\rangle\}. \qquad (4)$$

The main purpose of this paper is accounting SU(3)-symmetry breaking effects which lead to the difference between wave functions of Σ^*- and Ξ^*-hyperons from Δ.

The paper is organized as follows. In section 2 we present the main definitions and notations. In section 3 we investigate the properties of the leading twist WF $\varphi_{\Sigma^*}(x)$ and discuss the sum rules treatment procedure. In sections 4 and 5 we investigate the properties of the leading twist WF $T_{\Sigma^*}(x)$ and leading twist WF properties of the Ξ^*-hyperon respectively. In section 6 we present the asymptotic values of various hyperon formfactors. And in section 7 we discuss the results.

2. MAIN DEFINITIONS

The Σ^{*+}-hyperon leading twist wave functions are determined by the matrix element of the three-local operator:

$$\langle 0|\,\varepsilon^{ijk}\,u^i_\alpha(z_1)\,u^j_\beta(z_2)\,s^k_\gamma(z_3)\,|\Sigma^{*+}(p,\lambda=1/2)\rangle =$$

$$= \frac{f_{\Sigma^*}}{4\sqrt{3}}\{(pC)_{\alpha\beta}\,\Sigma^*_\gamma\,V_{\Sigma^*}(z_ip) + (p\gamma_5C)_{\alpha\beta}\,(\gamma_5\Sigma^*)_\gamma\,A_{\Sigma^*}(z_ip)\} +$$

$$+ \frac{f^T_{\Sigma^*}}{8\sqrt{3}}(\sigma_{\mu\nu}\,Cp_\nu)_{\alpha\beta}\,(\gamma_\mu\Sigma^*)_\gamma\,T_{\Sigma^*}(z_ip). \qquad (5)$$

Here: i, j and k are colour indicies; α, β and γ are spinor indicies; $|\Sigma^*(p)\rangle$ is the Σ^*-state with the momentum p; the constants f_{Σ^*} and $f^T_{\Sigma^*}$ determines the values of V_{Σ^*} and T_{Σ^*} at the origin; C is the charge conjugation matrix. The WF $V_{\Sigma^*}(x_1, x_2, x_3)$ (and similary $A_{\Sigma^*}(x)$ and $T_{\Sigma^*}(x)$)

$$V_{\Sigma^*}(z_i\,p) = \int_0^1 d_3x\,\exp[-i\Sigma\,x_i\,(z_ip)]\;V_{\Sigma^*}(x) \qquad (6)$$

describes the distribution of three quarks in $\Sigma^*(\lambda=1/2)$-state in longitudinal momentum fractions. The wave functions $V_{\Sigma^*}(x)$, $A_{\Sigma^*}(x)$ and $T_{\Sigma^*}(x)$ have the properties:

$$V_{\Sigma^*}(x_1,x_2,x_3) = V_{\Sigma^*}(x_2,x_1,x_3),$$
$$T_{\Sigma^*}(x_1,x_2,x_3) = T_{\Sigma^*}(x_2,x_1,x_3),$$
$$A_{\Sigma^*}(x_1,x_2,x_3) = -A_{\Sigma^*}(x_2,x_1,x_3). \qquad (7)$$

In the SU(3)-symmetry limit $f_\Delta = f_{\Sigma^*} = f^T_{\Sigma^*}$ and $T_\Delta(x_1, x_2, x_3) = T_{\Sigma^*}(x_1, x_2, x_3) = \varphi_{\Sigma^*}(x_1, x_3, x_2)$ But with SU(3)-symmetry breaking effects taken into account $T_{\Sigma^*}(x)$ can not be expressed through $V_{\Sigma^*}(x)$ and $A_{\Sigma^*}(x)$. The formula (5) is equivalent to the following form of the Σ^{*+}-state:

$$|\Sigma^{*+\uparrow}(\lambda=1/2)\rangle = \frac{1}{\sqrt{3}}\int_0^1 \frac{d_3x}{4\sqrt{24x_1x_2x_3}}\times$$
$$\times\{f_{\Sigma^*}[V(x)-A(x)]_{\Sigma^*}\,|u^\uparrow(x_1)\,u^\downarrow(x_2)\,s^\uparrow(x_3)\rangle +$$
$$+ f_{\Sigma^*}[V(x)+A(x)]_{\Sigma^*}\,|u^\downarrow(x_1)\,u^\uparrow(x_2)\,s^\uparrow(x_3)\rangle +$$
$$+ f^T_{\Sigma^*}\,T_{\Sigma^*}(x)\,|u^\uparrow(x_1)\,u^\uparrow(x_2)\,s^\downarrow(x_3)\rangle\}. \qquad (8)$$

The isotopic invariance determines the form of others Σ^*-states:

$$|\Sigma^{*-\uparrow}(\lambda=1/2)\rangle = \frac{1}{\sqrt{3}}\int_0^1 \frac{d_3x}{4\sqrt{24x_1x_2x_3}}\times$$
$$\times\{f_{\Sigma^*}[V(x)-A(x)]_{\Sigma^*}\,|d^\uparrow(x_1)\,d^\downarrow(x_2)\,s^\uparrow(x_3)\rangle +$$
$$+ f_{\Sigma^*}[V(x)+A(x)]_{\Sigma^*}\,|d^\downarrow(x_1)\,d^\uparrow(x_2)\,s^\uparrow(x_3)\rangle +$$
$$+ f^T_{\Sigma^*}\,T_{\Sigma^*}(x)\,|d^\uparrow(x_1)\,d^\uparrow(x_2)\,s^\downarrow(x_3)\rangle\}; \qquad (9)$$

$$|\Sigma^{*0\uparrow}(\lambda=1/2)\rangle = \frac{1}{\sqrt{3}}\int_0^1 \frac{d_3x}{4\sqrt{24x_1x_2x_3}}\times$$

$$\times\Big\{f_{\Sigma^*}[V(x)-A(x)]_{\Sigma^*}\,\Big|\frac{u^\uparrow(x_1)\,d^\downarrow(x_2)+d^\uparrow(x_1)\,u^\downarrow(x_2)}{\sqrt{2}}\,s^\uparrow(x_3)\Big\rangle +$$

$$+ f_{\Sigma^*}[V(x)+A(x)]_{\Sigma^*}\,\Big|\frac{u^\downarrow(x_1)\,d^\uparrow(x_2)+d^\downarrow(x_1)\,u^\uparrow(x_2)}{\sqrt{2}}\,s^\uparrow(x_3)\Big\rangle +$$

$$+ f^T_{\Sigma^*}\,T_{\Sigma^*}(x)\,\Big|\frac{u^\uparrow(x_1)\,d^\uparrow(x_2)+d^\uparrow(x_1)\,u^\uparrow(x_2)}{\sqrt{2}}\,s^\downarrow(x_3)\Big\rangle\Big\} \qquad (10)$$

3. WF $\varphi_{\Sigma^*}(x)$ OF Σ^*-HYPERON

To find the values of moments (n_1, n_2, n_3) for $\varphi_{\Sigma^*}(x)$ we use the following correlators:

$$I_{\Sigma^*}^{(n,000)}(q,z)=i\int dx\exp(iqx)\,\langle 0|T\{J_\tau^{(n)}(x)\,\bar{O}_{\tau'}^{(000)}(0)\}\,|0\rangle\times$$
$$\times\hat{z}_{\tau'\tau}=(zq)^{n_1+n_2+n_3+3}\,I_{\Sigma^*}^{(n,000)}(q^2)\,, \qquad (11)$$
$$(n)=(n_1,n_2,n_3)\,;$$
$$J_\tau^{(n)}(x)=\varepsilon^{ijk}[\,D^{n_1}u(x)]^i\,C\hat{z}(1-\gamma_5)\,[\,D^{n_2}u(x)]^j(\gamma_\mu(1+\gamma_5)\,[\,D^{n_3}s(x)]^k)_\tau\,,$$
$$\bar{O}_{\tau'}^{(000)}(0)=\varepsilon^{ijk}\{\bar{u}_{\tau'}^i(0)\,\bar{s}^k(0)\,\gamma_\mu\hat{z}C\bar{u}^j(0)+$$
$$+\bar{s}_{\tau'}^u(0)\,\bar{u}^k(0)\,\gamma_\mu\hat{z}C\bar{u}^j(0)+\bar{u}_{\tau'}^i(0)\,\bar{u}^k(0)\,\gamma_\mu\hat{z}C\bar{s}^j(0)\}\,,$$
$$D=z_\mu(i\partial_\mu-gA_\mu)\,,\qquad z^2=0\,. \qquad (12)$$

Note that because the current $\bar{O}_{\tau'}^{(000)}(0)$ is totaly symmetric in quarks interchange, the $\Sigma(1189)$-hyperon contribution into the correlator $I_{\Sigma^*}^{(n,000)}$ is zero in the exact SU(3)-symmetry limit. But, taking into account the SU(3)-symmetry breaking effects, we have to choose the spectral density in the form:

$$\frac{1}{\pi}\mathrm{Im}\,I_{\Sigma^*}^{(n,000)}(s)=\frac{4}{3}R_{\Sigma^*}^{(n)}\,\delta(s-M_{\Sigma^*}^2)+8R_\Sigma^{(n)}\,\delta(s-M_\Sigma^2)+\frac{\alpha_1^{(n)}s}{40\pi^4}\Theta(s-s^{(n)})\,,$$
$$R_{\Sigma^*}^{(n)}=f_{\Sigma^*}[\,2f_{\Sigma^*}+f_{\Sigma^*}^T]\,\varphi_{\Sigma^*}^{(n)}\,,$$
$$R_\Sigma^{(n)}=f_\Sigma[\,f_\Sigma-f_\Sigma^T]\,\varphi_\Sigma^{(n)}\,, \qquad (13)$$

where $R_\Sigma^{(n)}$ is the $\Sigma(1189)$-hyperon contribution (the Σ-hyperon leading twist WF was investigated in [5].), $R_{\Sigma^*}^{(n)}$ is the Σ^*-hyperon contribution and the last term in (13) is the perturbation theory contribution. The quantities s_n determine the beginning of the smooth continuum in correlators. The sum rules obtained from (11) and (12) have the form:

$$\frac{4}{3}R_{\Sigma^*}^{(n)}\exp(-M_{\Sigma^*}^2/M^2)+8R_\Sigma^{(n)}\exp(-M_\Sigma^2/M^2)=$$
$$=\frac{\alpha_1^{(n)}}{40\pi^4}M^4[1-(1+H)\exp(-H)]-\frac{\alpha_2^{(n)}m_s^2M^2}{8\pi^2}[1-\exp(-H)]+$$
$$+\frac{\alpha_3^{(n)}}{72\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle-\frac{16\alpha_4^{(n)}}{81\pi M^2}\langle\sqrt{\alpha_s}\,\bar{u}u\rangle^2+\frac{\alpha_5^{(n)}m_s}{3\pi^2}\langle\bar{s}s\rangle-$$
$$-\frac{\alpha_6^{(n)}m_s}{6\pi^2M^2}\langle\bar{s}ig\sigma_{\mu\nu}G_{\mu\nu}^a\frac{\lambda^a}{2}s\rangle,\qquad H=s^{(n)}/M^2. \qquad (14)$$

Here (and below) the hyperon masses are considered as known and are taken from the Particle Data Tables. In (14) the term $\alpha_1^{(n)}$ is

Fig. 1.

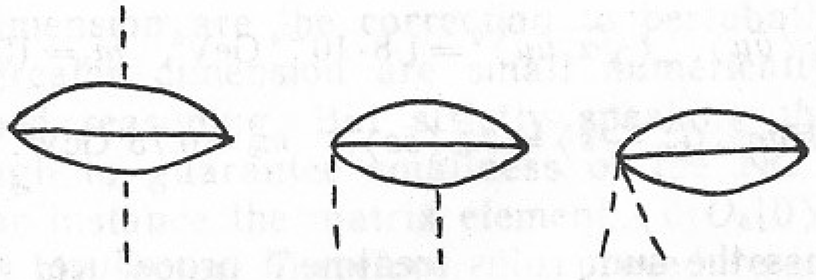

Fig. 2.

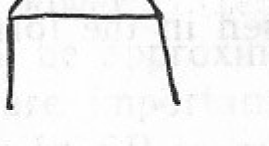

Fig. 3.

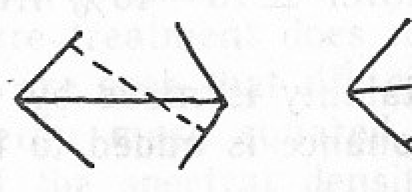

Fig. 4.

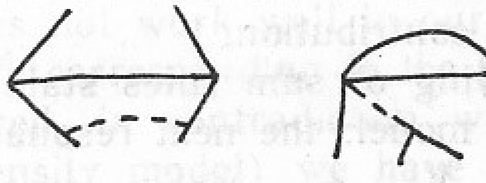

Fig. 5.

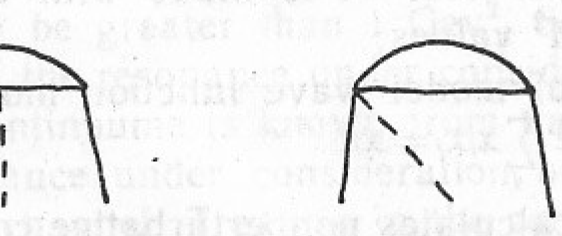

Fig. 6.

due to Fig. 1 diagram with massless quarks; the term $\alpha_2^{(n)}$ is a correction to this diagram due to the s-quark mass; the term $\alpha_3^{(n)}$ is due to Fig. 2 diagrams; the term $\alpha_5^{(n)}$ is due to Fig. 3 diagrams; the term $\alpha_4^{(n)}$ is due to Fig. 4 and Fig. 5 diagrams (these diagrams give both $\langle 0|\bar{u}u|0\rangle^2$ and $\langle 0|\bar{s}s|0\rangle^2$ condensates, and to simplify the formulae we have written the summary contribution using $\langle 0|\bar{s}s|0\rangle\simeq 0.8\times\langle 0|\bar{u}u|0\rangle$, (see [6]); the term $\alpha_6^{(n)}$ is due to diagrams. The coefficients $\alpha_i^{(n)}$ are given in Table 1. When fitting the sum rules we used the standard values of condensates (see, for instance, [6]):

$$\langle \frac{\alpha_s}{\pi} G^2 \rangle = 1.2 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad \langle \bar{u}u \rangle = -(0.25\ \mathrm{GeV})^3,$$

$$\langle \bar{s}s \rangle = 0.8 \langle \bar{u}u \rangle, \quad \langle \sqrt{\alpha_s}\,\bar{u}u \rangle^2 = 1.8 \cdot 10^{-4}\,\mathrm{GeV}^6, \quad m_s = 150\ \mathrm{MeV},$$

$$\langle \bar{s} i g \sigma_{\mu\nu} G^a_{\mu\nu} \frac{\lambda^a}{2} s \rangle = m_0^2 \langle \bar{s}s \rangle, \quad m_0^2 = 0.75\ \mathrm{GeV}^2.$$

Let us discuss the sum rules treatment procedure, which is used in this paper, in more detail.

Notice that the standart sum rules treatment procedure, which we used before [7, 8, 5], includes the following points:

1. The spectral density is chosen in the form: $\frac{1}{\pi}\mathrm{Im}\, I(s) =$ (resonance) + continuum;

2. The best fit is made in an interval of M^2 values where nonperturbative corrections are at the level $\simeq 10-40\%$ from the perturbation theory contribution;

3. The verifying of sum rules stability is made by varying the spectral density model: the next resonance is added to the spectral density, so that $\frac{1}{\pi}\mathrm{Im}\, I(s) =$ (resonance) + (next resonance) + continuum, and then the best fit is made with two resonances in the same interval of M^2 values;

4. The choice of model wave function moments is made, using the constraints like $\sum x_i x_j = x_i$.

As a rule one calculates nonperturbative corrections (below NC) only within a lowest dimension (in our case 4 and 6). Certainly, there are NC with a greater dimension (for instance, $\langle 0 | G^3 | 0 \rangle$ and others) at the SR, but usually it is implied that such contributions are smaller than the NC of a lowest dimension on the following grounds. First, we do the so-called «borelezation procedure», which suppresses NC with greater dimension by the factorial multiplier. Second, it is expected, that NC $\langle 0 | O_k | 0 \rangle$ with the dimenson $(\mathrm{mass})^{2k}$ is proportional to $(\mu_0)^{2k}$, where μ_0 is the characteristic scale, $\simeq$ (300—400) Mev. Since this matrix element enters into the

sum rules in combination $\langle 0 | O_k | 0 \rangle / M^{2k}$, then expansion parameter $(\mu_0)^2/M^2$ is small. At the same time, due to the special choice of the fitting interval the NC of the lowest dimension are less than 40% from the pertubation theory contribution. So, the usual SR ideology is: the main contribution into the SR is the perturbative one, NC of the lowest dimension are the correction to pertubative theory, and NC of the greater dimension are small numerically according to abovementioned reasoning. But strictly speaking these arguments are not enough to guarantee smallness of the NC of the greater dimension: for instance the matrix element $\langle 0 | O_k | 0 \rangle$ can be multiplied by very big factor. Therefore, this procedure is based heavily on the following hypothesis. Such scale of M_0^2 exists that power corrections hierarchy is valid above M_0^2 (that is NC of greater dimension give a smaller contribution into the SR than that of lower dimension), and the power corrections series blows up below M_0^2: so, this series can not be approximated by few first terms, since all the series members are important. If we take this hypothesis for granted than we may fit SR in any interval of M^2 (which is above M_0^2) —even where the contribution of lower dimension NC is greater than the perturbative theory contribution. Since the standard SR procedure treatment does not work well in our case (the most of SR (14) are such that $s^{(n)}$, corresponding to the best fit, is below the Σ^*-hyperon mass squared, in contradiction with the primary choosing of the spectral density model) we have to treat all SR (14) *fixing the fitting interval and fixing the begining of the smooth continuum* $s^{(n)}$. It is known from the work experience with the baryonic SR that M^2 must be greater than 1 Gev^2. So we can treat SR hereabout the mass of the resonance under consideration. The beginning of the smooth continuume is known from the physical reasons: it is above the resonance under consideration but belove the next resonance in a given channel. Treating in the common way the SR, for which the best fit results don't contradict the primary choise of the spectral density model, we'll find the begining of the smooth continuum $s^{(n)}$. It is the (001) moment in our case. The rest of the SR (14) we will treat with that $s^{(n)}$. Really this procedure terns out viable. Actually, the treating of the octet baryon correlators in this way gives the results close to the standard treating procedure results. So the hypothesis about the existence of the scale M_0^2 doesn't contradict to SR. The results which follow from fitting the sum rules (14) are presented in Table 2. In order to get some information about stability of these results we have calculated independent correlator $I_{\Sigma^*}^{(n,100)}(q, z)$ with the auxiliary current

$$\bar{O}_{\tau'}^{(100)}(0)=\varepsilon^{ijk}\{\bar{u}_{\tau'}^{i}(0)\,\bar{s}^{k}(0)\,\gamma_\mu\hat{z}C[\bar{u}(0)\,\overleftarrow{D}]^{j}+\bar{s}_{\tau'}^{i}(0)\times$$
$$\times\bar{u}^{k}(0)\,\gamma_\mu\hat{z}C[\bar{u}(0)\,\overleftarrow{D}]^{j}+\bar{u}_{\tau'}^{i}(0)\,\bar{u}^{k}(0)\,\gamma_\mu\hat{z}C[\bar{s}(0)\,\overleftarrow{D}]^{j}\}. \quad (15)$$

In order not to overload the text, we don't write here the explicit form of these sum rules. We present only the results which follow from fitting these sum rules.

It is seen from Table 2 that moment values of the WF $\varphi_{\Sigma^*}(x)$, obtained from two independent correlators, agree well with each other. Notice that moment values of the WF $\varphi_N(x)$, obtained from three independent proton correlators (000), (100) and (200), agree well with each other too [8]. It is necessary to emphasize that from our point of view this good agreement is neither trivial fact nor the accidental one. The reason is the following. The perturbation theory contribution into the values of WF moments is essentially different for different proton correlators: (000), (100) and (200) (e.g. for the moment (100) it is 1/3, 3/7 and 1/2 correspondingly). The NC contribution into the same moments of WF differs greatly too. But the change of the NC contribution from correlator to correlator is in such correspondence with the change of perturbation theory contribution that values of the WF moments practically do not change. This way NC «know» about the relative contribution of perturbation theory into the sum rules and correct the former to the «rignt» WF moment value. Besides, one more effect comes in to play in the decuplet case. The Σ-hyperon contribution into the correlator (000) is negligible since SU(3)-symmetry breaking effects for the octet residues are small. The corresponding contribution into the correlator (100) is not small because SU(3)-symmetry breaking effects for the shape of octet wave functions are not small. But in spite of that both correlators: (000) and (100) give the values of WF moments which agree well with each other.

The values of model wave function

$$\varphi_{\Sigma^*}(x)=\varphi_{as}(x)\,[-0.0927x_1^2+0.5295x_2^2-0.3735x_3^2+$$
$$+0.2364(x_1-x_2)+0.5042x_3-0.1533]$$

are presented in Table 2 also.

4. WF T_{Σ^*}

To find the $T_{\Sigma^*}(x)$ we used the following correlator:

$$K_{\Sigma^*}^{(n,000)}(q,z)=i\int dx\,\exp(iqx)\,\langle 0|\,T\{T_{\tau}^{(n)}(x)\,\bar{O}_{\tau'}^{(000)}(0)\}\,|0\rangle\times$$
$$\times\hat{z}_{\tau'\tau}=(zq)^{n_1+n_2+n_3+3}\,K_{\Sigma^*}^{(n,000)}(q^2)\,;$$
$$T_{\tau}^{(n)}(x)=\varepsilon^{ijk}[D^{n_1}u(x)]^{i}\,C\gamma_\mu\hat{z}(1+\gamma_5)\,[D^{n_2}u(x)]^{j}\,.$$
$$\{(1-\gamma_5)\,[D^{n_3}s(x)]^{k}\}_{\tau}\,. \quad (17)$$

The current $\bar{O}_{\tau'}^{(000)}(0)$ is the same as in (12). The spectral dencity is chosen in the form

$$\frac{1}{\pi}\,\mathrm{Im}\,K_{\Sigma^*}^{(n,000)}(s)=\frac{4}{3}S_{\Sigma^*}^{(n)}\,\delta(s-M_{\Sigma^*}^2)+16S_{\Sigma}^{(n)}\,\delta(s-M_{\Sigma}^2)+$$
$$+\frac{\beta_1^{(n)}s}{40\pi^4}\,\Theta(s-s^{(n)})\,,$$
$$S_{\Sigma^*}^{(n)}=f_{\Sigma^*}^{T}[2f_{\Sigma^*}+f_{\Sigma^*}^{T}]\;T_{\Sigma^*}^{(n)}\,, \quad (18)$$
$$S_{\Sigma}^{(n)}=f_{\Sigma}^{T}[f_{\Sigma}^{T}-f_{\Sigma}]\;T_{\Sigma}^{(n)}\,,$$

where $S_{\Sigma^*}^{(n)}$ is the $\Sigma^*(1189)$-hyperon contribution, $S_{\Sigma}^{(n)}$ is the Σ-hyperon contribution and the last term in (18) is the perturbation theory contribution.

The corresponding SR have the form

$$\frac{4}{3}\,S_{\Sigma^*}^{(n)}\exp(-M_{\Sigma^*}^2/M^2)+16S_{\Sigma}^{(n)}\exp(-M_{\Sigma}^2/M^2)=$$
$$=\frac{\beta_1^{(n)}}{40\pi^4}\,M^4[1-(1+H)\exp(-H)]-\frac{\beta_2^{(n)}m_s^2M^2}{8\pi^2}\,[1-\exp(-H)]+$$
$$+\frac{\beta_3^{(n)}}{72\pi^2}\langle\frac{\alpha_s}{\pi}\,G^2\rangle-\frac{16\beta_4^{(n)}}{81\pi M^2}\,\langle\sqrt{\alpha_s}\,\bar{u}u\rangle^2+\frac{\beta_5^{(n)}m_s}{3\pi^2}\,\langle\bar{s}s\rangle-$$
$$-\frac{\beta_6^{(n)}m_s}{\pi^2M^2}\,\langle\bar{s}ig\sigma_{\mu\nu}\,G_{\mu\nu}^{a}\,\frac{\lambda^a}{2}\,s\rangle. \quad (19)$$

The coefficients $\beta_i^{(n)}$ are given in Table 3. The results which follow from fitting the sum rules (19) are presented in Table 4. We have calculated also the independent correlator $K_{\Sigma^*}^{(n,100)}(q,z)$ with auxiliary current $\bar{O}_{\tau'}^{(100)}(0)$ (see (15)). The results which follow from fitting these sum rules are presented in Table 4 too. The moment values of the model WF

$$T_{\Sigma^*}(x)=\varphi_{ac}(x)\,[\,0.1647(x_1^2+x_2^2)-0.8205x_3^2+0.8594x_3-0.1925] \qquad (20)$$

are presented in Table 4 also. The constants f_{Σ^*} and $f^T_{\Sigma^*}$ are equal to

$$f_{\Sigma^*}\simeq f^T_{\Sigma^*}\simeq 1.5\cdot 10^{-2}\,\mathrm{GeV}^2,\qquad f_\Delta/f_{\Sigma^*}\simeq 1.25. \qquad (21)$$

Since in the exact SU(3)-symmetry limit $f_\Delta/f_{\Sigma^*}=1$, so it is seen from (21) that SU(3)-symmetry breaking effects are $\simeq 25\%$.

5. $\Xi^*(1530)$ HYPERON

The Ξ^*-hyperon leading twist wave functions are determined by the matrix element of three-local operator (see (5)) with the replacement $u\to s$, $s\to d$, $\Sigma^*\to\Xi^*$. In exact SU(3)-symmetry limit constant f_{Ξ^*} is equal to $f^T_{\Xi^*}$, and $\varphi_\Delta(x)=\varphi_{\Xi^*}(x)$.

Ξ^*-states can also be written in the form:

$$|\,\Xi^{*0\uparrow}(\lambda=1/2)\,\rangle=\frac{1}{\sqrt{3}}\int_0^1\frac{d_3x}{4\sqrt{24x_1x_2x_3}}\times$$

$$\times\{f_{\Xi^*}[\,V(x)-A(x)]_{\Xi^*}\,|\,s^\uparrow(x_1)\,s^\downarrow(x_2)\,u^\uparrow(x_3)\,\rangle+$$

$$+f_{\Xi^*}[\,V(x)+A(x)]_{\Xi^*}\,|\,s^\downarrow(x_1)\,s^\uparrow(x_2)\,u^\uparrow(x_3)\,\rangle+$$

$$+f^T_{\Xi^*}\,T_{\Xi^*}(x)\,|\,s^\uparrow(x_1)\,s^\uparrow(x_2)\,u^\downarrow(x_3)\,\rangle,$$

$$|\,\Xi^{*-\uparrow}(\lambda=1/2)\,\rangle=\frac{1}{\sqrt{3}}\int_0^1\frac{d_3x}{4\sqrt{24x_1x_2x_3}}\times$$

$$\times\{f_{\Xi^*}[\,V(x)-A(x)]_{\Xi^*}\,|\,s^\uparrow(x_1)\,s^\downarrow(x_2)\,d^\uparrow(x_3)\,\rangle+$$

$$+f_{\Xi^*}[\,V(x)+A(x)]_{\Xi^*}\,|\,s^\downarrow(x_1)\,s^\uparrow(x_2)\,d^\uparrow(x_3)\,\rangle+$$

$$+f^T_{\Xi^*}\,T_{\Xi^*}(x)\,|\,s^\uparrow(x_1)\,s^\uparrow(x_2)\,d^\downarrow(x_3)\,\rangle. \qquad (22)$$

For finding the WF $\varphi_{\Xi^*}(x)=V_{\Xi^*}(x)-A_{\Xi^*}(x)$ we have used correlators (11) and (12) with replacements $u\to s$, $s\to u$, $\Sigma^*\to\Xi^*$. The spectral density was chosen in the form (13) with the replacement $\alpha_i^{(n)}\to\gamma_i^{(n)}$. The corresponding sum rules have the form (14) with replacements $\Sigma\to\Xi$, $\alpha_i^{(n)}\to\gamma_i^{(n)}$. The coefficients $\gamma_i^{(n)}$ are given in Table 5. We have calculated also the independent correlator $I^{(n,100)}_{\Xi^*}(q,z)$ with auxiliary current $\bar{O}^{(100)}_{\gamma'}(0)$ (15) with the same replacements. The

results which follow from fitting these sum rules are presented in Table 6. The moment values of the model WF

$$\varphi_{\Xi^*}(x)=\varphi_{ac}(x)\,[\,-0.189x_1^2-0.015x_2^2+0.387x_3^2+$$

$$+0.09(x_1-x_2)-0.418x_3+0.137] \qquad (23)$$

are also given therein.

For finding the WF $T_{\Xi^*}(x)$ we have used correlators (17) with replacements $u\to s$, $s\to u$. The spectral dencity was chosen in the form (18) with the replacements $\Sigma^*\to\Xi^*$, $\Sigma\to\Xi$, $\beta_i^{(n)}\to\delta_i^{(n)}$. The corresponding sum rules have the form (19) with replacements $\beta_i^{(n)}\to\delta_i^{(n)}$, $\Sigma^*\to\Xi^*$, $\Sigma\to\Xi$. The coefficients $\delta_i^{(n)}$ are given in Table 7. We have calculated also the independent correlator $K^{(n,100)}_{\Xi^*}(q,z)$ with auxiliary current $\bar{O}^{(100)}_{\gamma'}(0)$ (15) with the same replacements. The results which follow from fitting these sum rules are presented in Table 8. The moment values of the model WF

$$T_{\Xi^*}(x)=\varphi_{ac}(x)\,[\,-0.147(x_1^2+x_2^2)+0.813x_3^2-0.844x_3+0.231] \qquad (24)$$

are also given therein. The constants f_{Ξ^*} and $f^T_{\Xi^*}$ are equal to

$$f_{\Xi^*}\simeq f^T_{\Xi^*}\simeq 1.5\cdot 10^{-2}\ \mathrm{GeV}^2,\qquad f_\Delta/f_{\Xi^*}\simeq 1.25. \qquad (25)$$

In the exact SU(3)-symmetry limit $f_\Delta=f_{\Xi^*}=f^T_{\Xi^*}$. It is seen from (25) that SU(3)-symmetry breaking effects are $\simeq 25\%$.

6. THE ELECTROMAGNETIC FORMFACTORS OF HYPERONS

These formfactors can be measured in the e^+e^--annihilation, $e^+e^-\to\bar{B}B$. It has been pointed out in [5, 9] that there is no need to do a new calculations of the Feynman diagrams for these formfactors, because one can use the results [10] for the nucleon formfactors (see below (28) and Table 9).

The properties of the leading twist wave function of the Ω^--hyperon have been investigated in paper [4] (see (2) and (3)). The value of the Ω^--hyperon formfactor is

$$Q^4F_1^{\Omega^-}(Q^2)\simeq -2\cdot 10^{-2}\ \mathrm{GeV}^4,\qquad I_{\Omega^-}\simeq -900. \qquad (26)$$

The values of the Δ-resonance formfactors were found in [9], those are:

$$Q^4F_1^{\Delta^{++}}(Q^2)\simeq 1.7\cdot 10^{-1}\ \mathrm{GeV}^4,$$

$$Q^4 F_1^{\Delta^+}(Q^2) \simeq 8.5 \cdot 10^{-2}\ \mathrm{GeV}^4,$$
$$|F_1^{\Delta^0}(Q^2)/F_1^{\Delta^+}(Q^2)| \ll 1,$$
$$Q^4 F_1^{\Delta^-}(Q^2) \simeq -8.5 \cdot 10^{-2}\ \mathrm{GeV}^4. \qquad (27)$$

The asymptotia of the nucleon formfactors has the form [10]:

$$Q^4 F_1(Q^2) = \frac{(4\pi\bar{\alpha}_s)^2}{54} |f_N|^2 I_N.$$

$$I_N = \int_0^1 d_3x\, d_3y \left\{ 2 \sum_{i=1}^{7} e_i T_i(x, y) + \sum_{i=8}^{14} e_i T_i(x, y) \right\}, \qquad (28)$$

where e_i is the charge of those quark which interacts with the photon in the given diagram, $e_d = e_s = -1/3$, $e_u = 2/3$. The explicit form of various diagrams and $T_i(x, y)$ are given in [10]. (For the reader convenience we present these results below in Table 9.)

Comparing the explicit forms of the proton and $\Sigma^{*+}(\lambda = 1/2)$-hyperon states:

$$|P\uparrow\rangle = f_N \int_0^1 \frac{d_3x}{4\sqrt{6}} \{\varphi_N(x)\, |u_1^\uparrow u_2^\downarrow d_3^\uparrow\rangle - T_N(x)\, |u_1^\uparrow u_2^\uparrow d_3^\downarrow\rangle,$$

$$|\Sigma^{*+}\uparrow\rangle = \frac{f_{\Sigma^*}}{\sqrt{3}} \int_0^1 \frac{d_3x}{4\sqrt{6}} \left\{ \varphi_{\Sigma^*}(x)\, |u_1^\uparrow u_2^\downarrow s_3^\uparrow\rangle + \frac{f_{\Sigma^*}^T}{2 f_{\Sigma^*}} T_{\Sigma^*}(x)\, |u_1^\uparrow u_2^\uparrow s_3^\downarrow\rangle, \right.$$

one can establish the correspondence:

$$f_N \to f_{\Sigma^*}/\sqrt{3}, \quad \varphi_N(x) \to \varphi_{\Sigma^*}(x),$$
$$T_N(x) \to -\frac{f_{\Sigma^*}^T}{2 f_{\Sigma^*}} T_{\Sigma^*}(x). \qquad (29)$$

Using (29) one can obtain the Σ^*-hyperon formfactors as follows.

1. Let us present the electromagnetic current in the form:

$$J_\mu^\gamma = 1/2\, J_\mu^v + 1/6\, J_\mu^0 - 1/3\, J_\mu^s,$$
$$J_\mu^v = \bar{u}\gamma_\mu u - \bar{d}\gamma_\mu d, \quad J_\mu^0 = \bar{u}\gamma_\mu u + \bar{d}\gamma_\mu d. \quad J_s^\mu = \bar{s}\gamma_\mu s.$$

2. Because there is no d-quarks in the Σ^{*+}-state, the formfactors of the currents J_μ^v and J_μ^0 coincide. The formfactor of the current J_μ^v

is obtained from (28) by the replacement (29) and with $e_u = 1$, $e_s = 0$. Using the explicit form of $T_i(x, y)$ from Table 9 and (16), (20) and (21), one obtains (here and below $\bar{\alpha}_s = 0.3$):

$$Q^4 F_{1\Sigma^*}^v(Q^2) = Q^4 F_{1\Sigma^*}^0(Q^2) \simeq 9.3 \cdot 10^{-2}\ \mathrm{GeV}^4,$$
$$I_{\Sigma^*}^v = I_{\Sigma^*}^0 \simeq 4.7 \cdot 10^3.$$

3. The formfactor of the current J_μ^s is obtained from (28) by the replacement (29) and with $e_u = 0$, $e_s = 1$ (i.e. accounting only for those diagrams NN 7, 12, 13 and 14 where the photon interacts with the s-quark). One obtains

$$Q^4 F_{1\Sigma^*}^s \simeq 0\ \mathrm{GeV}^4, \quad I_{\Sigma^*}^s \simeq 0.$$

The formfactors of Σ^*-hyperons have the form:

$$Q^4 F_1^{\Sigma^{*+}}(Q^2) = Q^4 \left[\frac{1}{2} F_{1\Sigma^*}^v + \frac{1}{6} F_{1\Sigma^*}^0 - \frac{1}{3} F_{1\Sigma^*}^s \right] \simeq 6.2 \cdot 10^{-2}\ \mathrm{GeV}^4,$$
$$Q^4 F_1^{\Sigma^{*0}}(Q^2) = Q^4 \left[\frac{1}{6} F_{1\Sigma^*}^0 - \frac{1}{3} F_{1\Sigma^*}^s \right] \simeq 1.6 \cdot 10^{-2}\ \mathrm{GeV}^4,$$
$$Q^4 F_1^{\Sigma^{*-}}(Q^2) = Q^4 \left[-\frac{1}{2} F_{1\Sigma^*}^v + \frac{1}{6} F_{1\Sigma^*}^0 - \frac{1}{3} F_{1\Sigma^*}^s \right] \simeq -3.1 \cdot 10^{-2}\ \mathrm{GeV}^4.$$

$$\frac{F_1^{\Delta^-}(Q^2)}{F_1^{\Omega^-}(Q^2)} \simeq 4.3; \quad \frac{F_1^{\Delta^+}(Q^2)}{F_1^{\Sigma^{*+}}(Q^2)} \simeq 1.4; \quad \frac{F_1^{\Delta^-}(Q^2)}{F_1^{\Sigma^{*-}}(Q^2)} \simeq 2.7. \qquad (30)$$

Let us point out that all these ratios are equal to unity in the SU(3)-limit. From (30) we see that though the value of SU(3)-symmetry breaking effects in the wave functions is $\simeq 20-30\%$, the SU(3)-symmetry breaking effects in the formfactors can be considerably greater.

Unfortinately the accuracy of the ratios (30) is not very good. The reason is the following. The magnetic formfactor value is very sensitive to the precise form of the WF when the properties of the latter are close to the properties of the asimptotic one. (Remind that the proton formfactor is zero when one take the asymptotic WF.) So really we choose the values of the Σ^*- and Ξ^*-hyperons WF moments inside the intervals allowed by sum rules so that all correlations, which follow from the exact SU(3)-symmetry limit, fulfil as good as possible. At the same time we have not succeeded in finding such model wave functions, which simultaneously satisfy sum rules requirements and lead to the small ($\simeq 20-30\%$) SU(3)-symmetry breaking effects in the formfactors.

The Ξ^{*0} and Ξ^{*-}-hyperons formfactors can be obtained in the similar way from (28) using the replacement $f_N \to f_{\Xi^*}/\sqrt{3}$, $\varphi_N(x) \to \varphi_{\Xi^*}(x)$, $T_N(x) \to -f_{\Xi^*}/2f_{\Xi^*} T_{c^*}(x)$. Using the explicit form (23) — (25) one obtains:

$$Q^4 F_1^{\Xi^{*0}}(Q^2) \simeq 1.4 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{\Xi^{*0}} \simeq 0.71 \cdot 10^3;$$

$$Q^4 F_1^{\Xi^{*-}}(Q^2) \simeq -8.3 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{\Xi^{*-}} \simeq -0.42 \cdot 10^4;$$

$$\frac{F_1^{\Delta^-}(Q^2)}{F_1^{\Xi^{*-}}(Q^2)} \simeq 1.0; \quad \frac{F_1^{\Sigma^{*0}}(Q^2)}{F_1^{\Xi^{*-}}(Q^2)} \simeq 1.1; \tag{31}$$

It is seen from (26), (27), (30) and (31) that typical values of the decuplet hyperons formfactors are $\simeq 10^{-2}$ Gev^4. The smallness of these formfactors (as compared with the nucleon formfactor which is equal to 0.95 GeV^4) is due to specific properties of the octet and decuplet wave functions. The first one is highly asymmetric, while the second one is close to symmetric. Therefore the typical values of the integrals including the decuplet wave functions are in order of madnitude smaller than the corresponding integrals for the octet wave functions.

The asymptotic behaviour of the transition formfactors $\gamma\Sigma^* \to \Sigma$ can be obtained as follows:

1. One has to replace $|f_N|^2 \to f_{\Sigma^*} f_\Sigma/\sqrt{3}$ in (28);

2. The expressions for $T_i(x, y)$ have the form (see [7, 10]) $T_i(x, y) = N_i(x, y)/D_i(x, y)$. One has to replace $1/D_i(x, y) = 1/2[1/D_i(x, y) + 1/D_i(y, x)]$.

3. In $N_i(x, y)$ one has to replace:

$$\varphi_N(x) \to \varphi_\Sigma(x), \quad T_N(x) \to \frac{-f_\Sigma^T}{2f_\Sigma} T_\Sigma(x), \quad \varphi_N(y) \to \varphi_{\Sigma^*}(y), \quad T_N(y) \to \frac{-f_{\Sigma^*}^T}{2f_{\Sigma^*}} T_{\Sigma^*}(y).$$

Using the explicit form (16), (20) and (21) one obtains:

$$Q^4 F_1^{\Sigma^{*+}\Sigma^+}(Q^2) = 3.3 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{\Sigma^{*+}\Sigma^+} \simeq 0.29 \cdot 10^4,$$

$$Q^4 F_1^{\Sigma^{*0}\Sigma^0}(Q^2) = 2.4 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{\Sigma^{*0}\Sigma^0} \simeq 0.21 \cdot 10^4,$$

$$Q^4 F_1^{\Sigma^{*-}\Sigma^-}(Q^2) = 1.6 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{\Sigma^{*-}\Sigma^-} \simeq 0.13 \cdot 10^4, \tag{32}$$

The transition formfactors $\gamma\Xi^* \to \Xi$ can be obtained in the similar way:

$$Q^4 F_1^{\Xi^{*-}\Xi^-}(Q^2) = 1.3 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{\Xi^{*-}\Xi^-} \simeq 1.1 \cdot 10^3,$$

$$Q^4 F_1^{\Xi^{*0}\Xi^0}(Q^2) = 2.4 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{\Xi^{*0}\Xi^0} \simeq 1.9 \cdot 10^3. \tag{33}$$

The transition formfactor $\gamma\bar{P}\Delta$ have been found in [9] and it's value is:

$$Q^4 F_1^{P\Delta^+}(Q^2) = 2 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{P\Delta^+} \simeq 2.1 \cdot 10^3. \tag{34}$$

The transition formfactor $\gamma\Sigma^{*0} \to \Lambda$ can be obtained by the analogous changes from the transition formfactor $\gamma\Sigma^0 \to \Lambda$ (see [5]):

$$Q^4 F_1^{\Sigma^0\Lambda} = \frac{f_\Lambda f_\Sigma}{\sqrt{2}} \frac{(4\pi\bar{\alpha}_s)^2}{54} I_{\Sigma^0\Lambda},$$

$$I_{\Sigma^0\Lambda} = \int d_3x\, d_3y \{2[T_1^{\Sigma^0\Lambda}(x, y) + T_9^{\Sigma^0\Lambda}(x, y)] - [(T_3^{\Sigma^0\Lambda}(x, y) + T_4^{\Sigma^0\Lambda}(x, y) + T_5^{\Sigma^0\Lambda}(x, y)) + (x \leftrightarrow y)]\},$$

$$Q^4 F_1^{\Sigma^{*0}\Lambda}(Q^2) = 1.5 \cdot 10^{-2}\,\mathrm{GeV}^4, \quad I_{\Sigma^{*0}\Lambda} \simeq 1.5 \cdot 10^3, \tag{35}$$

$$\frac{F_1^{\Delta^0 n}(Q^2)}{\sqrt{3/2}F_1^{\Sigma^{*0}\Lambda}(Q^2) - 1/2 F_1^{\Sigma^{*0}\Sigma^0}(Q^2)} \simeq 4.0; \quad \frac{F_1^{\Xi^{*0}\Xi^0}(Q^2)}{F_1^{\Delta^0 n}(Q^2)} \simeq 1.2. \tag{36}$$

It is seen from (32) — (35) that typical values of the octet-decuplet formfactors are $\simeq 10^{-2}$ Gev^4. The smallness of these formfactors (as compared with the nucleon formfactor $= 0.95$ GeV^4) is due to specific properties of the wave functions of hadrons participating in the process. The octet wave functions are highly asymmetric, while the decuplet one are close to symmetric. It has been pointed out in [11] that in such a situation there is a strong destructive interference of various diagram contributions which leads to a strong supression. It is necessary to note also that the following formfactors $F_{\Xi^*-\Xi^-}$, $F_{\Sigma^*-\Sigma^-}$ and $\sqrt{3}F_{\Sigma^{*0}\Sigma^0}(Q^2) + F_{\Sigma^{*0}\Lambda}(Q^2)$ are equal to zero in the exact SU(3)-symmetry limit. It is seen from (32) and (33) that values of these formfactors are the same order of magnitude as other formfactors are ($\sqrt{3}F_{\Sigma^{*0}\Sigma^0}(Q^2) + F_{\Sigma^{*0}\Lambda}(Q^2) = 5.7 \cdot 10^{-2}$ GeV^4). As a whole, it is seen from (26) — (37) that SU(3)-symmetry breaking effects in octet-decuplet and decuplet-decuplet formfactors can be $\simeq 100\%$.

7. DISCUSSION

The main purpose of this paper is to investigate the properties of the leading twist wave functions of the Σ^*- and Ξ^*-hyperons. The main results obtained in this paper are the following.

1. For a determination of values of six lowest independent WF moments of the Σ^*- and Ξ^*-hyperons all sum rules with $n_1+n_2+n_3<3$ which follow from the two independent correlators have been obtained and treated. It is checked that the treatment of all independent sum rules leads to the results which agree well with each other.

2. Based on the knowledge of six lowest independent moments with $n_1+n_2+n_3<3$ we have proposed model wave functions for the Σ^*- and Ξ^*-hyperons.

3. It is checked that sum rules for the octet baryons doesn't contradict to the hypothesis about existence of such scale M_0^2, above which the power corrections hierarchy is valid.

4. Comparing the residue values of the decuplet baryons (see (3), (21) and (25)) it is seen that SU(3)-symmetry breaking effects are $\simeq 25\%$. At the same time the corresponding residues of the octet barions are practically the same as in exact SU(3)-symmetry limit.

5. Comparing the WF of various baryons (see (1), (2), (16), (20), (23) and (24)) it is seen that SU(3)-symmetry breaking effects are $\simeq 20-30\%$ both in octet WF and in decuplet WF. Let us consider for instance the following components of wave functions:

$$|\Delta^{+\uparrow}\rangle \to |u^\uparrow(x_1)\, u^\uparrow(x_2)\, d^\uparrow(x_3)\rangle,$$
$$|\Sigma^{*+\uparrow}\rangle \to |u^\uparrow(x_1)\, u^\uparrow(x_2)\, s^\uparrow(x_3)\rangle,$$
$$|\Xi^{*-\uparrow}\rangle \to |s^\uparrow(x_1)\, s^\uparrow(x_2)\, d^\uparrow(x_3)\rangle,$$
$$|\Omega^{-\uparrow}\rangle \to |s^\uparrow(x_1)\, s^\uparrow(x_2)\, s^\uparrow(x_3)\rangle \qquad (37)$$

and compare the values of their first moments:

$$|\Delta^{+\uparrow}\rangle : \{\langle x_1\rangle : \langle x_2\rangle : \langle x_3\rangle \simeq (32{:}36{:}32)\,\%\},$$
$$|\Sigma^{*+\uparrow}\rangle : \{\langle x_1\rangle : \langle x_2\rangle : \langle x_3\rangle \simeq (30{:}29{:}41)\,\%\},$$
$$|\Xi^{*-\uparrow}\rangle : \{\langle x_1\rangle : \langle x_2\rangle : \langle x_3\rangle \simeq (42{:}32{:}26)\,\%\},$$
$$|\Omega^{-\uparrow}\rangle : \{\langle x_1\rangle : \langle x_2\rangle : \langle x_3\rangle \simeq (33{:}33{:}33)\,\%\}. \qquad (38)$$

Comparing (37) and (38) we see that when u-quark is replaced by s-quark the momentum fraction it carries increases. This is characteristic for the meson wave functions [6] and for the octet baryon wave functions [5] also.

At the same time SU(3)-symmetry breaking effects influence on the octet and decuplet formfactors by different ways. The typical SU(3)-symmetry breaking effects in octet-octet formfactors are $\simeq 10-30\%$:

$$\frac{F_1^{\Sigma^+}}{F_1^p} \simeq 1.25\,; \quad \frac{-F_1^{\Sigma^0}}{F_1^n} \simeq 1.21\,; \quad \frac{F_1^{\Xi^0}}{F_1^n} \simeq 1.15\,; \quad \frac{F_1^{\Xi^-}}{F_1^{\Sigma^-}} \simeq 0.92\,;$$
$$\frac{2F_1^{\Lambda}}{F_1^n} \simeq 1.03\,; \quad \frac{-F_1^{\Lambda}}{F_1^{\Sigma^0}} \simeq 0.85\,; \quad \frac{-F_1^{\Sigma\Lambda}}{\sqrt{3}\,F_1^{\Lambda}} \simeq 1.35\,. \qquad (39)$$

At the same time it seems that the SU(3)-symmetry breaking effects in the octet-decuplet and decuplet-decuplet formfactors are greater. For instance:

$$\frac{F_1^{\Delta^-}(Q^2)}{F_1^{\Omega^-}(Q^2)} \simeq 4.3\,; \quad \frac{F_1^{\Delta^+}(Q^2)}{F_1^{\Sigma^{*+}}(Q^2)} \simeq 1.4\,; \quad \frac{F_1^{\Delta^-}(Q^2)}{F_1^{\Sigma^{*-}}(Q^2)} \simeq 2.7\,;$$
$$\frac{F_1^{\Delta^0 n}(Q^2)}{\sqrt{3}/2F_1^{\Sigma^{*0}\Lambda}(Q^2) - 1/2F_1^{\Sigma^{*0}\Sigma^0}(Q^2)} \simeq 4.0\,. \qquad (40)$$

However the results (40) have considerably less accurasy then (39), because formfactor values of the decuplet baryons turn out very sensitive to the precise form of the wave functions.

At the same time we have not succeeded in finding such model wave functions which simultaneously fulfil the sum rules and lead to the small SU(3)-symmetry breaking effects in the formfactors.

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Table 1

Moments	$\alpha_1^{(n)}$	$\alpha_2^{(n)}$	$\alpha_3^{(n)}$	$\alpha_4^{(n)}$	$\alpha_5^{(n)}$	$\alpha_6^{(n)}$
(000)	1	1	1	6.60	1	1
(100)	1/3	2/5	1/2	2.36	1/2	1/2
(010)	1/3	2/5	0	1.88	1/2	5/6
(001)	1/3	1/5	1/2	2.36	0	−1/3
(200)	1/7	1/5	3/5	581/600	3/10	1/3
(020)	1/7	1/5	1/5	323/600	3/10	2/3
(002)	1/7	1/15	3/5	34/30	0	0
(110)	2/21	2/15	−1/10	113/150	1/5	1/3
(101)	2/21	1/15	0	383/600	0	−1/6
(011)	2/21	1/15	−1/10	353/600	0	−1/6

Table 2

$(V-A)_{\Sigma^*}$					
Moments	Correlator (000)		Correlator (100)		Model WF
	SR	«FIT»	SR	«FIT»	
(000)	1	1	1	1	1
(100)	0.21−0.39	0.30	0.25−0.44	0.34	0.2993
(010)	0.19−0.38	0.28	—	—	0.2870
(001)	0.32−0.51	0.42	0.33−0.53	0.43	0.4137
(200)	0.10−0.18	0.14	0.10−0.19	0.14	0.1146
(020)	0.09−0.17	0.13	—	—	0.1313
(002)	0.18−0.26	0.22	0.15−0.24	0.20	0.1887
(110)	0.03−0.09	0.06	0.04−0.09	0.06	0.0577
(101)	0.08−0.13	0.11	0.11−0.17	0.14	0.1270
(011)	0.07−0.12	0.10	0.07−0.12	0.09	0.0980

Table 3

Moments	$\beta_1^{(n)}$	$\beta_2^{(n)}$	$\beta_3^{(n)}$	$\beta_4^{(n)}$	$\beta_5^{(n)}$	$\beta_6^{(n)}$
(000)	1	1	1	7.56	1	1/3
(100)	1/3	2/5	1/2	2.74	1/2	41/900
(001)	1/3	1/5	0	2.08	0	109/450
(200)	1/7	1/5	3/5	1.53	3/10	11/150
(002)	1/7	1/15	1/5	1.10	0	2/9
(110)	2/21	2/15	0	0.72	1/5	−17/450
(101)	2/21	1/15	−1/10	0.49	0	1/100

Table 4

T_{Σ^*}					
Moments	Correlator (000)		Correlator (100)		Model WF
	SR	«FIT»	SR	«FIT»	
(000)	1	1	1	1	1
(100)	0.17 — 0.35	0.26	—	—	0.253
(001)	0.38 — 0.58	0.48	0.34 — 0.54	0.44	0.494
(200)	0.08 — 0.16	0.12	—	—	0.080
(002)	0.22 — 0.30	0.26	0.19 — 0.28	0.24	0.220
(110)	0.01 — 0.06	0.03	—	—	0.036
(101)	0.09 — 0.14	0.11	0.08 — 0.14	0.10	0.137

Table 5

Moments	$\gamma_1^{(n)}$	$\gamma_2^{(n)}$	$\gamma_3^{(n)}$	$\gamma_4^{(n)}$	$\gamma_5^{(n)}$	$\gamma_6^{(n)}$
(000)	1	2	1	4.56	2	1
(100)	1/3	3/5	1/2	1.30	1/2	−1/6
(010)	1/3	3/5	0	1.78	1/2	1/2
(001)	1/3	4/5	1/2	1.48	1	2/3
(200)	1/7	4/15	3/5	53/150	3/10	1/6
(020)	1/7	4/15	1/5	61/75	3/10	2/3
(002)	1/7	2/5	3/5	53/150	3/5	1/2
(110)	2/21	2/15	−1/10	59/150	0	−1/3
(101)	2/21	1/5	0	139/300	1/5	0
(011)	2/21	1/5	−1/10	43/75	1/5	1/6

Table 6

$(V-A)_{\Xi^*}$					
Moments	Correlator (000)		Correlator (100)		Model WF
	SR	«FIT»	SR	«FIT»	
(000)	1	1	1	1	1
(100)	0.26 — 0.50	0.41	0.21 — 0.57	0.39	0.417
(010)	0.25 — 0.49	0.38	0.22 — 0.51	0.36	0.318
(001)	0.14 — 0.38	0.22	0.08 — 0.43	0.25	0.265
(200)	0.13 — 0.23	0.19	0.06 — 0.22	0.14	0.201
(020)	0.10 — 0.21	0.16	0.07 — 0.17	0.12	0.134
(002)	0.05 — 0.15	0.08	0.00 — 0.13	0.05	0.111
(110)	0.09 — 0.16	0.14	0.10 — 0.19	0.14	0.123
(101)	0.05 — 0.12	0.09	0.05 — 0.16	0.11	0.093
(011)	0.04 — 0.11	0.07	0.06 — 0.14	0.10	0.061

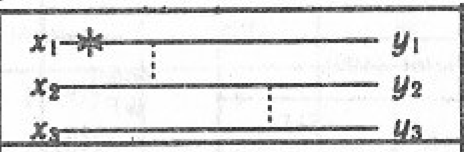
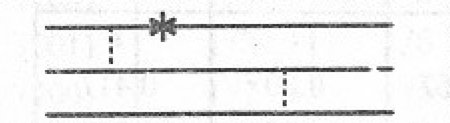
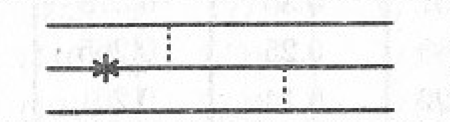
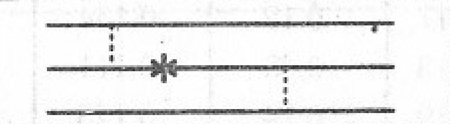
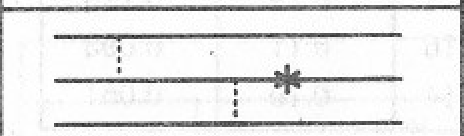
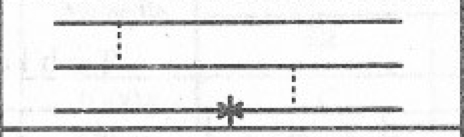
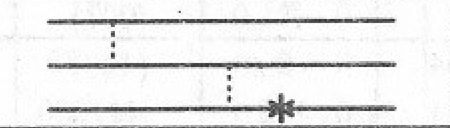
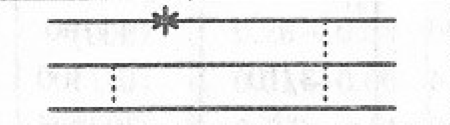
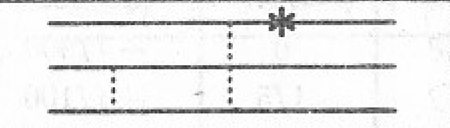
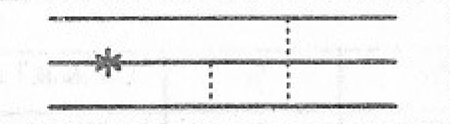
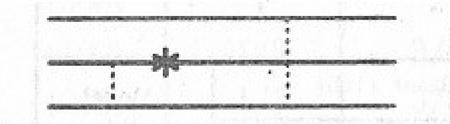
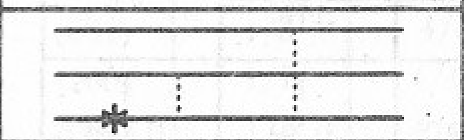
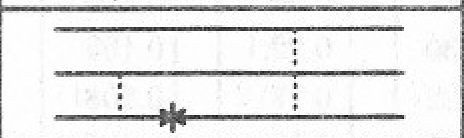
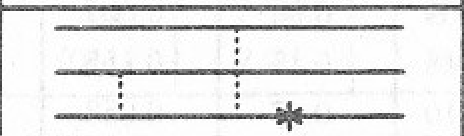
Table 7

Moments	$\delta_1^{(n)}$	$\delta_2^{(n)}$	$\delta_3^{(n)}$	$\delta_4^{(n)}$	$\delta_5^{(n)}$	$\delta_6^{(n)}$
(000)	1	2	1	6.84	2	16/75
(100)	1/3	3/5	1/2	2.44	1/2	41/900
(001)	1/3	4/5	0	1.96	1	11/90
(200)	1/7	4/15	3/5	1.44	3/10	19/300
(002)	1/7	2/5	1/5	0.92	3/5	32/225
(110)	2/21	2/15	0	0.48	0	−7/900
(011)	2/21	1/5	−1/10	0.52	1/5	−1/100

Table 8

T_{Ξ^*}					
Moments	Correlator (000)		Correlator (100)		Model WF
	SR	«FIT»	SR	«FIT»	
(000)	1	1	1	1	1
(100)	0.27 — 0.52	0.37	0.29 — 0.49	0.39	0.416
(001)	0.09 — 0.34	0.21	0.14 — 0.30	0.22	0.168
(200)	0.13 — 0.23	0.18	0.12 — 0.22	0.17	0.208
(002)	0.02 — 0.13	0.07	0.03 — 0.09	0.06	0.062
(110)	0.11 — 0.18	0.15	0.12 — 0.18	0.15	0.155
(101)	0.04 — 0.11	0.07	0.05 — 0.10	0.07	0.053

Table 9

i	Diagram	$T_i(x, y)$
1		$\frac{(V(x)-A(x))(\dot{V}(y)-\dot{A}(y))+4T(x)\dot{T}(y)}{(1-x_1)^2x_3(1-y_1)^2y_3}$
2		0
3		$\frac{-4T(x)\dot{T}(y)}{x_1(1-x_2)x_3y_1(1-y_1)y_3}$
4		$\frac{(V(x)-A(x))(\dot{V}(y)-\dot{A}(y))}{x_1x_3(1-x_3)y_1(1-y_1)y_3}$
5		$\frac{-(V(x)+A(x))(\dot{V}(y)+\dot{A}(y))}{x_1x_3(1-x_3)y_1(1-y_2)y_3}$
6		0
7		$\frac{(V(x)-A(x))(\dot{V}(y)-\dot{A}(y))+(V(x)+A(x))(\dot{V}(y)+\dot{A}(y))}{x_1(1-x_3)^2y_1(1-y_3)^2}$
8		0
9		$\frac{(V(x)-A(x))(\dot{V}(y)-\dot{A}(y))+4T(x)\dot{T}(y)}{(1-x_1)^2x_2(1-y_1)^2y_2}$
10		$\frac{(V(x)+A(x))(\dot{V}(y)+\dot{A}(y))+4T(x)\dot{T}(y)}{x_1(1-x_2)^2y_1(1-y_2)^2}$
11		0
12		$\frac{-(V(x)+A(x))(\dot{V}(y)+\dot{A}(y))}{x_1x_2(1-x_3)y_1y_2(1-y_2)}$
13		$\frac{4T(x)\dot{T}(y)}{x_1(1-x_1)x_2y_1y_2(1-y_2)}$
14		$\frac{-(V(x)-A(x))(\dot{V}(y)-\dot{A}(y))}{x_1(1-x_1)x_2y_1y_2(1-y_3)}$

A.A. Ogloblin

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А.А. Оглоблин

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