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RELATIVISTIC QUARK MODEL IN INFINITE MOMENTUM FRAME AND STATIC PROPERTIES OF HADRONS

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РЕЛЯТИВИСТСКАЯ МОДЕЛЬ КВАРКОВ В СИСТЕМЕ БЕСКОНЕЧНОГО ИМПУЛЬСА И СТАТИЧЕСКИЕ СВОЙСТВА АДРОНОВ

В системе бесконечного импульса сформулирована модель мезонов и барионов, построенных из релятивистских кварков. Построен вариант этой модели, допускающий классификацию по группе SU(6)xO(3). В этом варианте модели проведен анализ данных по магнитным моментам, электромагнитным радиусам и отношению G_A/G_V для нуклона и получено их самосогласованное описание. Получены оценки на среднеквадратичные импульсы кварков в нуклоне, которые оказались сравнимыми с их массой, на массу, и аномальные магнитные моменты u и d - кварков. С этими значениями параметров получены предсказания для амплитуд радиационных распадов резонансов P_{33}(I232), P_{11} (I470), P_{11}(I7I0) и резонансов, входящих в мультиплет [70,I$^-$].

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A model of mesons and baryons composed of relativistic quarks is formulated in the infinite momentum frame. A variant of this model allowing the classification over SU(6)xO(3) group is constructed. In this variant of the model the analysis of the data on the magnetic moments, electromagnetic radii and on the ratio G_A / G_V for nucleon is carried out and their selfconsistent description is obtained. Estimates for mass, anomalous magnetic moments and mean square momenta of quarks in nucleons are obtained. Using these parameters we predict the radiative decay amplitudes of P_{33} (1232), P_{11} (1470), P_{11} (1780) resonances and nucleon resonances of $[70, 1^-]$ multiplet which are in good agreement with experimental data.

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1. Introduction

It is well known that there are quite enough convincing arguments in favour of hadrons treated as bound states of constituent quarks. Here should be mentioned many relations between amplitudes and cross sections of exclusive reactions, some features of many-particle reactions and hadron-nuclear collisions at high energies (see, e.g., reviews [1, 2, 3]). The nonrelativistic quark model describes well the magnetic moments of the baryon octet ($\frac{1}{2}^{+}$) and in a natural way explains the relations between electromagnetic radii of nucleons [1, 2]. In this model also many radiative decays of mesons and nucleon resonances are well described (see, e.g., [2, 4])

At the same time, when obtaining the baryon resonance mass spectrum by nonrelativistic formulae (see, e.g., [5]), it turns out that quark mean square momentum in nucleon $\langle Q^2 \rangle$ (it is connected with the parameter α used in Ref. [5] by the relation $\langle Q^2 \rangle = \alpha^2$) is not small being comparable with its squared mass m_q^2 . When trying to obtain the numerical values of the nucleon electromagnetic radii by the nonrelativistic formulae, $\langle Q^2 \rangle / m_q^2$ turns out not small, too. The incorrectness of the nonrelativistic approximation is especially evident for the radiative decays of the meson and nucleon resonances, when the emitted photon momentum is comparable with the quark mass, and the quark interacting with photon turns out necessarily relativistic. The predictions of the quark model for these decays were made by the nonrelativistic formulae, but with relativistic parameters. Thus to clear up whether the quark model is really applicable to the description of the hadron static characteristics the correct calculations, taking into account the quark relativistic motion in hadrons, are necessary.

The relativistic-invariant model of mesons and baryons treated as bound states of relativistic quarks has been formulated in Refs.[6-9] in the light-front dynamics. In this model we analyzed magnetic moments of baryon octet ($\frac{1}{2}^{+}$) as well as radiative decays of vector mesons [10]. In particular, it was shown that these quantities can be described well at sufficiently large quark momentum.

The similar questions were investigated also by Leutwyller and Stern (see Ref.[11] , where one can find other references on this subject). In Ref.[12] it was shown on the example of the pion formfactor that the results of the relativistic model [6] can be derived more obviously by the consideration of the time ordered diagrams of the old-fashioned perturbation theory in the infinite momentum frame (IMF).The IMF (as well as the light-front dynamics) is used in order to minimize the vacuum fluctuations. As a result the space-time picture of the interaction turns out analogous to that in the nonrelativistic quantum mechanics. This makes it possible to define the hadron-quarks transitions vertex functions (VF) by analogy with the nonrelativistic bound state wave functions. In the diagram approach we found the general method of construction of the VF for many-particle bound states, corresponding to the variants of spins and orbital angular momenta composition considered in Refs. [6-9]. We have also found the variant of the VF construction which corresponds to the experimentally observed SU(6)xO(3) classification of hadron states.

In sec.4 we consider the role of the relativistic effects in the description of the magnetic moments, electromagnetic radii and ratio G_A / G_V for nucleon assuming that it is pure $[56, 0^+]$ state. A simultaneous analysis of the experimental data on these quantities is carried out and the main parameters of the model are determined: mass of the constituent u- and d-quarks, their mean square momenta and anomalous magnetic moments. The considered experimental data are described well by the set of parameters obtained.

In secs. 5-8 we consider the role of relativistic effects in the description of the radiative decays of nucleon resonances. The predictions for these decays are obtained which agree well with the experiment. Note, in particular, that in our approach: 1) $\Delta(1232) \to N\gamma$ decay amplitude agrees with the experiment better than in nonrelativistic quark model; 2) we have obtained the agreement with the experiment for signs of $N(1470) \to N\gamma$ decay amplitudes; 3) in our approach $A^p_{3/2}(D_{13}(1520)) > A^p_{1/2}(S_{11}(1535))$, which agrees well with the experiment and cannot be obtained in the nonrelativistic model.

2. Main kinematic relations

Consider the electromagnetic transition

$$B \to B' + \gamma^*, \tag{1}$$

which is described by the set of Fig.1 time ordered diagrams of the old-fashioned perturbation theory. For relativistic quarks together with Fig.1a diagram, generally speaking, also contribute the rest diagrams of Fig.1 which violate the space-time picture of the process (1), corresponding to the nonrelativistic quantum mechanics. However in specially chosen IMF ($P \to \infty$), in which $K_0 = -K_3 = \frac{M^2 - M'^2 - K_\perp^2}{4P}$ [12] (K is the virtual photon momentum, M and M' are masses of initial and final hadrons), the Fig.1b, c, ... diagrams give no contribution at least for longitudinal components of the electromagnetic current. For transverse components of the electromagnetic current the suppression of Fig.1b, c,... diagrams connected with energy denominators can be, in principle, compensated by a large contribution of the vertices, since for the current transverse components the vertex $\gamma q \to q$ is suppressed as compared with the vertex $\gamma \to q\bar{q}$ as $\frac{1}{P}$. Therefore the use of transverse components of electromagnetic current requires a detailed analysis of Fig.1b, c, ... diagrams contribution.

Let us write the main kinematic relations for the Fig.1a diagram in the IMF. In this section we consider the three-quark states only, the generalization for many-quark bound states can be done in a simple way. For the initial quark momenta (denote them a , b , c) we introduce the Feynman parametrization:

$$\vec{P}_i = x_i \vec{P} + \vec{K}_{i\perp}, \quad \sum_{i=a,b,c} \vec{K}_{i\perp} = 0, \quad \sum_{i=a,b,c} x_i = 1, \tag{2}$$

$$E_i = |x_i| P + \frac{m_i^2 + \vec{K}_{i\perp}^2}{2P|x_i|} \; .$$

In $P \to \infty$ system the main contribution to the matrix element of Fig.1a diagram gives the integration region, where

$$0 < x_i < 1. \tag{2'}$$

In what follows we shall assume that the photon interacts with quark c multiplying the obtained results by 3. For the final quark momenta we obtain:

$$\begin{gathered} \vec{P}_i' = \vec{P}_i \equiv x_i \vec{P}' + \vec{K}_{i\perp}', \quad \vec{K}_{i\perp}' = \vec{K}_{i\perp} + x_i \vec{K}_\perp, \quad i = a, b, \\ \vec{P}_c' = \vec{P}_c - \vec{K}_\perp \equiv x_c \vec{P}' + \vec{K}_{c\perp}', \quad \vec{K}_{c\perp}' = \vec{K}_{c\perp} - (1 - x_c) \vec{K}_\perp , \end{gathered} \tag{3}$$

where $\vec{P}' = \vec{P} - \vec{K}_\perp$ is the momentum of the final hadron.

The invariant masses of initial and final quark systems are equal to:

$$M_0^2 = \sum_{i=a,b,c} \frac{\vec{K}_{i\perp}^2 + m_i^2}{x_i} \; , \tag{4}$$

$$M_0'^2 = \sum_{i=a,b,c} \frac{\vec{K}_{i\perp}'^2 + m_i^2}{x_i} = M_0^2 + \frac{-2\vec{K}_{c\perp}\vec{K}_\perp + \vec{K}_\perp^2(1-x_c)}{x_c} .$$

It is convenient to introduce the variables K_{iz} :

$$K_{iz} + \omega_i = M_0 x_i , \tag{5}$$

$$\omega_i = \sqrt{\vec{K}_{i\perp}^2 + K_{iz}^2 + m_i^2} , \qquad M_0 = \sum_{i=a,b,c} \omega_i .$$

Variables $K_i\left(\vec{K}_{i\perp}, K_{iz}, \omega_i\right)$ have the meaning of the 4-momentum quarks in the c.m.s. of initial quarks. Recall that we consider diagrams of the noncovariant perturbation theory in the system $P \to \infty$, hence the energy does not conserve in the vertices $B(B') \to 3q$ and consequently the c.m.s. of the quarks does not coincide with the hadron rest frame, and the invariant masses M_0^2 and $M_0'^2$ do not coincide with hadron masses M^2 and M'^2. Formulae (5) can be obtained by the Lorentz transformation along the third axis (we assume that this axis coincides with $\vec{P}$) from the c.m.s. of the quarks to IMF.

In an analogous way we define the final quark momenta in the $\vec{P}' = 0$ frame:

$$K_{iz}' + \omega_i' = M_0' x_i , \tag{5'}$$

$$\omega_i' = \sqrt{\vec{K}_i'^2 + m_i^2} , \qquad M_0' = \sum_{i=a,b,c} \omega_i' .$$

Let us give the connection of the above-introduced variables with ξ , η $\vec{q}$, $\vec{Q}$ that are used in Refs.[7-9] :

$$\begin{aligned} x_a &= \xi\eta , & \vec{K}_{a\perp} &= \vec{q}_\perp + \xi\vec{Q} , \\ x_b &= (1-\xi)\eta , & \vec{K}_{b\perp} &= -\vec{q}_\perp + (1-\xi)\vec{Q}_\perp , \\ x_c &= 1-\eta , & \vec{K}_{c\perp} &= -\vec{Q}_\perp , \end{aligned} \tag{6}$$

$$q_z + \varepsilon_a = M_{ab}\,\xi\,,\quad M_{ab} = \varepsilon_a + \varepsilon_b\,,\quad \varepsilon_a = \sqrt{\vec{q}^{\,2} + m_a^2}\,,\quad \varepsilon_b = \sqrt{\vec{q}^{\,2} + m_b^2}\,, \tag{7}$$

$$Q_z + E_{ab} = M_0\,\eta\,,\quad E_{ab} = \sqrt{M_{ab}^2 + \vec{Q}^2}\,,$$

$$E_c = \sqrt{\vec{Q}^2 + m_c^2}\,,\qquad M_0 = E_c + E_{ab}\,.$$

The corresponding variables ξ', η', $\vec{q}_\perp'$, $\vec{Q}_\perp'$ for the final quarks are connected with ξ, η, $\vec{q}_\perp$, $\vec{Q}_\perp$ by the relations

$$\xi' = \xi\,,\quad \eta' = \eta\,,\quad \vec{q}_\perp{}' = \vec{q}_\perp\,,\quad \vec{Q}_\perp' = \vec{Q}_\perp + \eta\vec{K}_\perp\,. \tag{8}$$

In the variables ξ, η, $\vec{q}_\perp$, $\vec{Q}_\perp$ and $\vec{Q}_\perp'$ the invariant masses M_0^2, $M_0'^2$, M_{ab}^2 and $M_{ab}'^2$ have the form:

$$M_0^2 = \frac{\vec{Q}_\perp^2}{\eta(1-\eta)} + \frac{M_{ab}^2}{\eta} + \frac{m_c^2}{1-\eta}\,, \tag{9}$$

$$M_{ab}^2 = \frac{\vec{q}_\perp^{\,2}}{\xi(1-\xi)} + \frac{m_a^2}{\xi} + \frac{m_b^2}{1-\xi} = M_{ab}'^2\,,$$

$$M_0'^2 = \frac{\vec{Q}_\perp'^2}{\eta(1-\eta)} + \frac{M_{ab}^2}{\eta} + \frac{m_c^2}{1-\eta}\,.$$

The transition (1) matrix element corresponding to Fig.1a diagram has the form

$$\langle P', S' | J_\mu | P, S \rangle = \frac{3}{(2\pi)^6} \int \frac{d\vec{p}_a\, d\vec{p}_b}{16 E_a E_b E_c E_c'}\ \frac{1}{(E - E_a - E_b - E_c)} \tag{10}$$

$$\cdot\ \frac{1}{(E' - E_a - E_b - E_c')} \sum_{\substack{s_a, s_b,\\ s_c, s_c'}} \left[\Gamma^{S'}_{s_a, s_b, s_c'}(p_a, p_b, p_c')\right]^{+} \bar{u}_{s_c'}(p_c')\Big(Q_c\gamma_\mu$$

$$+ \frac{\varkappa_c}{2m_c}\,\sigma_{\mu\nu} K_\nu\Big) u_{s_c}(p_c)\quad \Gamma^{S}_{s_a, s_b, s_c}(p_a, p_b, p_c)\,,$$

where S are the particle spin projections, Q_c, $\frac{\varkappa_c}{2m_c}$, m_c are respectively the charge, anomalous magnetic moment (a.m.m.) and mass of quark c, $\sigma_{\mu\nu} = \frac{1}{2}(\gamma_\mu\gamma_\nu - \gamma_\nu\gamma_\mu)$, Γ are the vertex functions corresponding to $B \to 3q$ transition. In the $P \to \infty$ frame the energy denominators in (10) are equal to

$$E - E_a - E_b - E_c = \frac{M^2 - M_0^2}{2P} , \qquad (11)$$

$$E' - E_a - E_b - E_c' = \frac{M'^2 - M_0'^2}{2P} .$$

Introducing the notations

$$\Psi^{S}_{s_a, s_b, s_c}(P_a, P_b, P_c) = \frac{\Gamma^{S}_{s_a, s_b, s_c}(P_a, P_b, P_c)}{M^2 - M_0^2} , \qquad (12)$$

$$\Psi^{S'}_{s_a, s_b, s_c'}(P_a, P_b, P_c') = \frac{\Gamma^{S'}_{s_a, s_b, s_c'}(P_a, P_b, P_c')}{M'^2 - M_0'^2} ,$$

we rewrite (10) for longitudinal components of the electromagnetic current in $P \to \infty$ frame in the form:

$$\frac{1}{2P} \langle P', S' | J_{0,3} | P, S \rangle = 3 \sum_{\substack{s_a, s_b, \\ s_c, s_c'}} \int d\Gamma \left[\Psi^{S'}_{s_a, s_b, s_c'}(P_a, P_b, P_c') \right]^+ \qquad (13)$$

$$\omega^+_{s_c'} \left(Q_c + i K_i \varepsilon_{ik} \sigma_k \frac{\varkappa_c}{2 m_c} \right) \omega_{s_c} \Psi^{S}_{s_a, s_b, s_c}(P_a, P_b, P_c) ,$$

where we use the normalization condition $\bar{u} u = 2m$, $d\Gamma$ is the phase space which can be presented in one of the following forms [7]:

$$4 \cdot (2\pi)^6 d\Gamma = \frac{d\vec{k}_{a\perp} d\vec{k}_{b\perp} dx_a dx_b}{x_a x_b x_c} = \frac{M_0 d\vec{k}_a d\vec{k}_b}{\omega_a \omega_b \omega_c} = \qquad (14)$$

$$= \frac{d\vec{q}_\perp d\vec{Q}_\perp d\xi \, d\eta}{\eta(1-\eta)\xi(1-\xi)} = \frac{M_{ab}}{\varepsilon_a \varepsilon_b} \frac{M_0}{E_c E_{ab}} d\vec{Q} d\vec{q} .$$

The normalization condition of vertex functions Ψ follows from the normalization of the elastic formfactor $B \to B\gamma$ at $P = P'$ and $S = S'$:

$$\int \sum_{S_a, S_b, S_c} \left| \Psi^{S}_{S_a, S_b, S_c} (P_a, P_b, P_c) \right|^2 d\Gamma = 1. \tag{15}$$

For the type (1) processes (as distinct, for example, from the quantities of the type $\langle \rho | J_\mu | 0 \rangle$) one may not write quark color indices since the color factors are similar in (13) and (15) and do not contribute to the final results. Vertex functions for baryons must be completely symmetric relative to quark replacements, so we write them in the form:

$$\Psi^{S}_{S_a, S_b, S_c} (P_a, P_b, P_c) = \mathcal{Y}(P_a, P_b, P_c) \psi^{S}_{S_a, S_b, S_c} (P_a, P_b, P_c), \tag{16}$$

where ψ is the spin-orbital part of the vertex functions, and the radial part of the vertex function $\mathcal{Y}(P_a, P_b, P_c)$ describes the momentum distribution of quarks in baryons being symmetric relative to their replacements. In the present work we shall assume that the function $\mathcal{Y}$ depends only on one variable, i.e. on the invariant mass of the quark system. When getting numerical results we use for the ground state the following form of this function:

$$\mathcal{Y}(P_a, P_b, P_c) \equiv \phi(M_0^2) = N \exp\left(-\frac{M_0^2}{6\alpha^2}\right), \tag{17}$$

which in nonrelativistic approximation turns into the oscillator wave function. Parameter N is determined from the normalization condition (15).

The detailed knowledge of the interaction dynamics being unknown, we cannot construct unambiguously the radial parts of the higher orbital excitations vertex functions corresponding to the ground state wave functions (17). In this work in sec.7, 8 we construct these functions by analogy with nonrelativistic oscillator model.

In the further calculations it will be convenient to use the matrix elements (13) written in the variables $\xi, \eta, \vec{q}_\perp, \vec{Q}_\perp, \vec{Q}'_\perp = \vec{Q}_\perp + \eta \vec{K}_\perp$:

$$\frac{1}{2P}\langle P',S'|J_{0,3}|P,S\rangle = 3\cdot\sum_{\substack{s_a,s_b,\\ s_c,s_c'}}\int d\Gamma\Big[\Psi^{S'}_{s_a,s_b,s_c'}(\xi,\eta,\vec{q}_\perp,\vec{Q}_\perp+\eta\vec{K}_\perp)\Big]^{+}\omega^{+}_{s_c'}\Big(Q_c+iK_i\varepsilon_{ik}\sigma_k\frac{\mathcal{H}_c}{2m_c}\Big)\omega_{s_c}\Psi^{S}_{s_a,s_b,s_c}(\xi,\eta,\vec{q}_\perp,\vec{Q}_\perp). \tag{13'}$$

<u>3. Construction of spin-orbital parts of vertex functions</u>

A. Two-body bound states

When constructing the vertex functions for two-body states we require that they must be written in the Lorentz-covariant form, so that spin and orbital angular momenta composition rule hold in the c.m.s. of constituent particles. Since the angular momenta composition law in the c.m.s. of two particles has the same form as in the nonrelativistic case[9], the vertex functions of two-particle states in c.m.s. coincide with the corresponding nonrelativistic wave functions. To write down these functions in the relativistic-covariant form for particles with spin 1/2 we introduce the bispinors

$$\mathcal{U}_s(P_i,P_0)=P^{+}(P_i,P_0)\mathcal{U}_s(P_i)\,,\quad \mathcal{V}_{-s}(-P_i,-P_0)=P^{-}(P_i,P_0)\mathcal{V}_{-s}(-P_i)\,, \tag{18}$$

where $P^{\pm}$ are projection operators which are equal to

$$P^{\pm}(P_i\,,\,P_0)=\frac{M_0\pm\hat{P}_0}{\sqrt{2(p_iP_0+m_iM_0)}}\,, \tag{19}$$

$\vec{P}_0=\vec{P}$, $E_0=\sqrt{M_0^2+\vec{P}_0^2}$, $P_{i0}=\sqrt{\vec{P}_i^2+m_i^2}$. Bispinors (18) satisfy Dirac equations $(\hat{P}_0-M_0)\mathcal{U}_s(p_i\,,P_0)=0$ and $(\hat{P}_0+M_0)\mathcal{V}_{-s}(-p_i,-P_0)=0$ and differ from bispinors $\mathcal{U}_s(P_0)$ and $\mathcal{V}_{-s}(-P_0)$ obtained by the Lorentz transformation from the $\vec{P}_0=0$ system by spin rotation:

$$v_s(P_i, P_o) = \sum_{s'} u_{s'}(P_o) V_{s's}(P_i, P_o), \quad V_{s's}(P_i, P_o) = \frac{\bar{u}_{s'}(P_o) u_s(P_i)}{\sqrt{2(P_i P_o + m_i M_o)}}, \tag{20}$$

$$v_{-s}(-P_i, -P_o) = \sum_{s'} v_{-s'}(-P_o)\left[V_{s's}(P_i, P_o)\right]^*. \tag{21}$$

In the $\vec{P}_o = 0$ system

$$V_{s's}(P_i, P_o)\Big|_{\vec{P}_o=0} = \delta_{ss'} . \tag{22}$$

In the $P \to \infty$ system the matrix

$$V_{s's}(P_i, P_o) \equiv U_{s's}(K_i) = \left[\frac{m_i + \omega_i + K_{iz} + i\,\varepsilon_{\ell m}\sigma_\ell K_{im}}{\sqrt{\vec{K}_{i\perp}^2 + (m_i + \omega_i + K_{iz})^2}}\right]_{s's} \tag{23}$$

coincides with Melosh matrix which was obtained in the current algebra as the transformation from the current quarks to constituent ones [13]. In the c.m.s. of constituent particles bispinor $u_s(P_i, P_o)$ has no lower component and coincides with $u_s(P_i)$ in $\vec{P}_i = 0$ frame:

$$\frac{1}{\sqrt{2M}} u_s(P_i, P_o)\Big|_{\vec{P}_o=0} = \frac{1}{\sqrt{2m_i}} u_s(P_i)\Big|_{\vec{P}_i=0} = \begin{pmatrix} w_s \\ 0 \end{pmatrix}. \tag{24}$$

Analogously bispinor $v_{-s}(-P_i, -P_o)$ in $\vec{P}_o = 0$ frame has no upper component. Introduce also four-dimensional vectors

$$\widetilde{P}_{i\mu} = P_{i\mu} - \frac{P_i P_o}{M_o^2} P_{o\mu}, \tag{25}$$

which in the c.m.s. of constituent particles $\vec{P}_o = 0$ have the form

$$\widetilde{P}_{i4} = 0, \quad \vec{\widetilde{P}}_i = \vec{P}_i = \vec{K}_i, \tag{26}$$

and tensor

$$\tilde{g}_{\mu\nu} = g_{\mu\nu} - \frac{P_{0\mu} P_{0\nu}}{M_0^2} , \qquad \tilde{g}_{\mu\nu} P_{0\mu} = \tilde{g}_{\mu\nu} P_{0\nu} = 0, \tag{27}$$

for which in the $\vec{P}_0 = 0$ frame we have

$$\tilde{g}_{0\mu} = \tilde{g}_{\mu 0} = 0 , \quad \tilde{g}_{ij} = g_{ij} , \qquad i,j = 1,2,3 \tag{28}$$

In the terms of vectors $\tilde{P}_{i\mu}$ and tensor $\tilde{g}_{\mu\nu}$ the "nonrelativistic" kinematics takes place in $\vec{P}_0 = 0$ frame , in other words , every possible products of the four vectors reduce to the products of their space components. Therefore, using $\tilde{P}_{i\mu}$, $\tilde{g}_{\mu\nu}$ and bispinors (18) one can easily write nonrelativistic vertex functions of two-particle system in the relativistic-covariant form. As an example we write down the vertex functions of π -meson and ρ -meson with $L = 0$ and $L = 2$ (S and D waves in ρ -meson):

$$\Psi^{\pi}_{s_1, s_2} = \frac{i}{\sqrt{2}} \frac{1}{2M_0} \bar{u}_{s_1}(p_1, P_0) \gamma_5 v_{-s_2}(-p_2, -P_0) \varphi^{\pi} , \tag{29}$$

$$\Psi^{\rho}_{s, s_1, s_2}(L = 0) = \frac{1}{2\sqrt{2} M_0} \left[\varphi^{\rho}_{\mu}(P_0)\right]_s \bar{u}_{s_1}(p_1, P_0) \gamma_{\mu} v_{-s_2}(-p_2, -P_0), \tag{30}$$

$$\Psi^{\rho}_{s, s_1, s_2}(L = 2) = -\frac{1}{2\sqrt{2} M_0} \left[\varphi^{\rho}_{\mu}(P_0)\right]_s \bar{u}_{s_1}(p_1, P_0) \gamma_{\nu} v_{-s_2}(-p_2, -P_0). \tag{31}$$

$$\left(\tilde{q}_{\mu} \tilde{q}_{\nu} - \frac{1}{3} \tilde{g}_{\mu\nu} \tilde{q}^2\right),$$

where $\tilde{q}_{\mu} = \tilde{P}_{1\mu} - \tilde{P}_{2\mu}$. In the $\vec{P}_0 = 0$ system the vertex functions (29, 30 , 31) have the nonrelativistic form:

$$\Psi^{\pi}_{s_1, s_2}\Big|_{\vec{P}_0 = 0} = \frac{i}{\sqrt{2}} \omega^{+}_{s_1} \sigma_2 \omega_{s_2} \varphi^{\pi} , \tag{29'}$$

$$\Psi^{\rho}_{s,s_1,s_2}(L=0)\Big|_{\vec{P}_0=0} = \frac{1}{\sqrt{2}}\, e_i^{s}\, \omega^{+}_{s_1}\, \sigma_i\, \sigma_2\, \omega_{s_2}\,, \qquad (30')$$

$$\Psi^{\rho}_{s,s_1,s_2}(L=2)\Big|_{\vec{P}_0=0} = \frac{1}{\sqrt{2}}\, e_i^{s}\, \omega^{+}_{s_1}\, \sigma_j\, \sigma_2\, \omega_{s_2}\left(q_i q_j - \frac{\delta_{ij}}{3}\vec{q}^{\,2}\right), \qquad (31')$$

where e_i^{s} are polarization vectors corresponding to $s = \pm 1.0$ for ρ - meson. In the $\rho \to \infty$ system the vertex functions (29, 30, 31) written through spinors

$$W_{s_i} \equiv U_{s_i' s_i}(K_i)\, \omega_{s_i'}\,, \qquad (32)$$

have the same form as in c.m.s. (29' - 31') and are equivalent to the wave functions constructed in Refs. [6, 9] in the light-front dynamics.

To obtain analogous results for particles with arbitrary spins one must together with projection operators (19) introduce the analogous operator for vector wave functions:

$$P_{\mu\nu}(P_a, P_0) = g_{\mu\nu} - \frac{(m_a P_{0\mu} + M_0 P_{a\mu}) P_{0\nu}}{M_0 (P_a P_0 + m_a M_0)}\,, \quad (P_0)_\mu P_{\mu\nu}(P_a, P_0) = 0. \qquad (33)$$

From the vector wave function $\varphi^{s}_{\mu}(P_a)$ let us constructed the function

$$\varphi^{s}_{\mu}(P_a, P_0) = P_{\mu\nu}(P_a, P_0)\, \varphi^{s}_{\nu}(P_a)\,. \qquad (34)$$

In the $\vec{P}_0 = 0$ system vector $\varphi^{s}_{\mu}(P_a, P_0)$ has no time component, and its space components coincide with space components of $\varphi^{s}_{\mu}(P_a)$ in the system $\vec{P}_a = 0$:

$$\varphi^{s}_{0}(P_a, P_0)\Big|_{\vec{P}_0=0} = 0, \quad \varphi^{s}_{i}(P_a, P_0)\Big|_{\vec{P}_0=0} = \varphi^{s}_{i}(P_a)\Big|_{\vec{P}_a=0} = e^{s}_{i}\,,\ i=1,2,3. \qquad (35)$$

Vector $\varphi^{s}_{\mu}(P_a, P_0)$ differs from $\varphi^{s}_{\mu}(P_0)$ by spin rotation:

$$\varphi^{i}_{\mu}(p_a, P_0) = V^{ji}(p_a, P_0)\,\varphi^{j}_{\mu}(P_0), \qquad i,j = 1,2,3, \tag{36}$$

where

$$V^{ji}(p_a, P_0) = -\varphi^{i}_{\mu}(p_a)\,\varphi^{j}_{\mu}(P_0) + \frac{[\varphi^{i}_{\mu}(p_a)P_{0\mu}][\varphi^{j}_{\nu}(P_0)p_{a\nu}]}{p_a P_0 + m_a M_0}. \tag{37}$$

In the $\vec{P}_0 = 0$ system

$$V^{ij} = \delta_{ij}. \tag{38}$$

In the $P \to \infty$ system

$$V^{ij}(p_a, P_0)\Big|_{P\to\infty} = \delta_{ij} - \frac{(\vec{K}_{a\perp})_i(\vec{K}_{a\perp})_j + \vec{K}^2_{a\perp}\delta_{i3}\delta_{j3} + (\omega_a + m_a + K_{az})[\delta_{i3}(\vec{K}_{a\perp})_j - \delta_{j3}(\vec{K}_{a\perp})_i]}{(\omega_a + m_a)(\omega_a + K_{az})} \tag{39}$$

Formula (39) coincides with the Melosh matrix in the vector representation.

B. Three-particle bound states

In Ref.[9] three-particle states are constructed in the following way: First, the state of two particles (a, b) with a definite total angular momentum and its projection is to be constructed. Then, considering the system (a, b) as an elementary particle, the state (a, b)c with the given value of total angular momentum is constructed according to the two-particle state construction rule. However it turns out that the three-quark states constructed in such a way cannot reproduce the experimentally observed SU(6)xO(3) hadron resonance spectrum [14]. Let us illustrate this on the example of three-quark bound states with $\ell_{ab} = L = 0$ (ℓ_{ab} is orbital angular momentum in (a, b) system, L is orbital angular momentum of c-quark relative to this system) and with proton quantum numbers.

Let us write down the corresponding vertex functions making use of the results of the previous section. In the case of $S_{ab} = T_{ab} = 0$ (S_{ab} and T_{ab} are spin and isospin of (a, b) quark system) we obtain:

$$\left[(\psi_1)^{s}_{s_a,s_b,s_c}\right]^{+} = \frac{1}{8M_{ab}M_0}\,\bar{u}_s(P_0)\,u^{u}_{s_c}(p_c,P_0)\ . \tag{40}$$

$$\left[\tilde{u}^{u}_{s_a}(p_a,P_{ab})C\gamma_5\,u^{d}_{s_b}(p_a,P_{ab}) - \tilde{u}^{d}_{s_a}(p_a,P_{ab})C\gamma_5\,u^{u}_{s_b}(p_b,P_{ab})\right],$$

where the upper indices stand for quark flavors; $P_{ab} = p_a + p_b$. The expression in square brackets corresponds to the bound state of quark pair (a, b) with $S_{ab} = T_{ab} = 0$. In the case of $S_{ab} = T_{ab} = 1$ we obtain

$$\left[(\psi_2)^{s}_{s_a,s_b,s_c}\right]^{+} = \frac{i}{4\sqrt{2}\,M_0M_{ab}}\left\{\sqrt{\frac{2}{3}}\,\bar{u}_s(P_0)\gamma_5\gamma_\mu u^{d}_{s_c}(p_c,P_0)P_{\mu\nu}(P_{ab},P_0)\cdot\right.$$

$$\cdot\,\tilde{u}^{u}_{s_a}(p_a,P_{ab})C\gamma_\nu\,u^{u}_{s_b}(p_b,P_{ab}) - \frac{1}{\sqrt{3}}\,\bar{u}_s(P_0)\gamma_5\gamma_\mu u^{u}_{s_c}(p_c,P_0)P_{\mu\nu}(P_{ab},P_0) \tag{41}$$

$$\left.\cdot\,\frac{1}{\sqrt{2}}\left[\tilde{u}^{u}_{s_a}(p_a,P_{ab})C\gamma_\nu\,u^{d}_{s_b}(p_b,P_{ab}) + u^{d}_{s_a}(p_a,P_{ab})C\gamma_\nu\,u^{u}_{s_b}(p_b,P_{ab})\right]\right\}.$$

Here the projection operator $P_{\mu\nu}(P_{ab}, P_0)$ is used for quarks (a, b) as they are in the state with $S_{ab} = 1$. Let us write (40, 41) in the system $P \to \infty$. Using the relation

$$P_{\mu\nu}(P_{ab},P_0)\,\tilde{u}(p_a,P_{ab})\,C\gamma_\nu\,u(p_b,P_{ab}) = \tag{42}$$

$$= \frac{M_{ab}}{M_0}\,\tilde{u}(p_a,P_{ab})\,\tilde{P}^{+}(P_{ab},P_0)C\gamma_\nu P^{+}(P_{ab},P_0)\,u(p_b,P_{ab}),$$

which can be obtained with the help of (19) and (33) and the relation

$$P^{+}(P_{ab},P_{0})P^{+}(P_i,P_{ab})\mathcal{U}_s(P_i)=\mathcal{U}_{s''}(P_0)V_{s''s'}(P_{ab},P_0)V_{s's}(P_i,P_{ab})\rightarrow \tag{43}$$

$$\xrightarrow[P\to\infty]{} \mathcal{U}_{s''}(P_0)\,\mathcal{U}_{s''s'}(\bar{P}_{ab})\,\mathcal{U}_{s's}(\bar{K}_i)\,, \qquad i=a,b,$$

where

$$\bar{P}_{ab}=(\vec{Q},\,E_{ab}),\quad \bar{K}_a=(\vec{q},\varepsilon_a),\qquad \bar{K}_b=(-\vec{q},\,\varepsilon_b), \tag{43'}$$

$\mathcal{U}_{s''s'}(\bar{P}_{ab})$ and $\mathcal{U}_{s's}(\bar{K}_i)$ are Melosh matrices defined according to (23). In the $P\to\infty$ system we obtain

$$\left[(\psi_1)^{S}_{s_a,s_b,s_c}\right]^{+}=\frac{i}{2}\,\mathcal{U}_{s_c s_c'}(K_c)\,\mathcal{U}_{s_b s_b'}(\bar{K}_b)\,\mathcal{U}_{s_a s_a'}(\bar{K}_a)\cdot \tag{44}$$

$$\cdot\,\omega_S^{+}\,\omega^{u}_{s_c'}\left(\tilde{\omega}^{u}_{s_a'}\sigma_2\,\omega^{d}_{s_b'}-\tilde{\omega}^{d}_{s_a'}\sigma_2\,\omega^{u}_{s_b'}\right),$$

$$\left[(\psi_2)^{S}_{s_a,s_b,s_c}\right]^{+}=\frac{i}{\sqrt{2}}\,\mathcal{U}_{s_c s_c'}(K_c)\,\mathcal{U}_{s_b s_b'}(\bar{P}_{ab})\,\mathcal{U}_{s_b' s_b''}(\bar{K}_b)\cdot \tag{45}$$

$$\cdot\,\mathcal{U}_{s_a s_a'}(\bar{P}_{ab})\,\mathcal{U}_{s_a' s_a''}(\bar{K}_a)\Big[\sqrt{\tfrac{2}{3}}\,\omega_S^{+}\sigma_\ell\,\omega^{d}_{s_c'}\left(\tilde{\omega}^{u}_{s_a''}\sigma_\ell\,\omega^{u}_{s_b''}\right)-$$

$$-\frac{1}{\sqrt{6}}\,\omega_S^{+}\sigma_\ell\,\omega^{u}_{s_c'}\left(\tilde{\omega}^{u}_{s_a''}\sigma_\ell\,\omega^{d}_{s_b''}+\tilde{\omega}^{d}_{s_a''}\sigma_\ell\,\omega^{u}_{s_b''}\right)\Big].$$

The vertex functions (44, 45) are equivalent to the corresponding wave functions constructed in Ref.[9] in the light-front dynamics.

In order to obtain the vertex functions for the three-quark bound state which corresponds to a proton, the vertex functions (40, 44) and (41, 45) must be symmetrized relative to the quark replacements. In the nonrelativistic limit the obtained vertex functions coincide with each other being equivalent to the proton $[56, 0^+]$ wave function. In the relativistic case these vertex functions differ from each other. For example, their quadratic over momenta terms include in the quark c.m.s. $\vec{P}_o = 0$ the following multiplets:

$$\psi_1: \quad [56, 0^+]_\tau, \; [20, 0^+], \tag{46}$$

$$\psi_2: \quad [56, 0^+]_\tau, \; [70, 0^+], \; [70, 2^+].$$

Thus the vertex functions ψ_1 and ψ_2 are not equivalent being not orthogonal. To extract the orthogonal spectrum it is necessary to have the whole series of states with the highest orbital angular momenta. Therefore we do not see the way the orthogonal basis which could be set in correspondence with the observed spectrum of baryon resonances can be built up on the basis of three-quark states with the definite values of quantum numbers $S_{a\ell}$, $\ell_{a\ell}$ and L.

2. In the present work, when constructing the three-quark states, we shall proceed from the requirement that they could be set in correspondence with the SU(6)xO(3) baryon resonance spectrum. Therefore we assume that in the quark c.m.s. $\vec{P}_o = 0$ the vertex functions written through two-component spinors have SU(6)xO(3) nonrelativistic structure. These vertex functions can be written in relativistic covariant form making use of bispinors $u(P_i, P_o)$ (18), vectors $\tilde{P}_{i\mu}$ (25) and tensor $\tilde{g}_{\mu\nu}$ (27). One can show that in the $P \to \infty$ system these vertex functions expressed through the spinors W_{S_i} (32) have the same form as in c.m.s. of quarks.

Let us write down, as an example, the vertex functions for proton and $\Delta^{++}(S_{31})$ resonance assuming that they are respectively pure $[56, 0^+]$ and $[70, 1^-]$ states:

$$\left[(\psi_p)^S_{S_a, S_b, S_c}\right]^+ = \frac{1}{4\sqrt{2}M_o^2}\bar{u}_S(P_o)\,u^u_{S_c}(P_c, P_o)\left[\tilde{u}^u_{S_a}(P_a, P_o)C\gamma_5\, u^d_{S_b}(P_b, P_o) - \right. \tag{47}$$

$$-\tilde{u}^{d}_{s_a}(p_a,P_0)C\gamma_5 u^{u}_{s_b}(p_b,P_0)\big] + \text{transmutations of } (a,b,c),$$

$$\left[(\psi_{\Delta^{++}})^{s}_{s_a,s_b,s_c}\right]^{+} = (2M_0)^{-2}\bar{u}_s(P_0)\,u^{u}_{s_c}(P_c,P_0)\cdot$$

$$\left[\tilde{u}^{u}_{s_a}(p_a,P_0)\,C\gamma_\mu\,\tilde{g}_{\mu\nu}\,\tilde{a}_\nu\,u^{u}_{s_b}(p_b,P_0)\right] + \text{transmutations of } (a,b,c) \qquad (48)$$

where

$$a_\nu = P_{a\nu} + P_{b\nu} - 2P_{c\nu}\,. \qquad (49)$$

In the $P\to\infty$ system we obtain

$$\left[(\psi_p)^{s}_{s_a,s_b,s_c}\right]^{+} = \frac{i}{\sqrt{2}}\,\omega^{+}_{s}W^{u}_{s_c}\left(\tilde{W}^{u}_{s_a}\sigma_2 W^{d}_{s_b} - \tilde{W}^{d}_{s_a}\sigma_2 W^{u}_{s_b}\right) + (a,b,c), \qquad (50)$$

$$\left[(\psi_{\Delta^{++}})^{s}_{s_a,s_b,s_c}\right]^{+} = \omega^{+}_{s}W^{u}_{s_c}\left(\tilde{W}^{u}_{s_a}\sigma_2\vec{\sigma}\,\vec{a}\,W^{u}_{s_b}\right) + (a,b,c). \qquad (51)$$

It is worth noting that for such three-particle vertex functions both the spin and orbital angular momentum of the quark pair in its rest frame have no definite value. For example, if we write the expression in the square brackets in (47) and (48) in the rest frame of (a, b) quark pair, we shall not get the simple expression without relative momentum of these quarks as it is the case in the vertex functions (40) and (41). Nevertheless, when constructing and classifying the states we shall use the quantum numbers corresponding to the nonrelativistic wave functions. Note that in this case the vertex functions for states with $L=0$, in particular for nucleon, are equivalent to the wave functions of Ref.[7].

4. Magnetic moments, charge and magnetic radii and ratio G_A/G_V for nucleons

In this section we shall consider the role of the relativistic effects in the description of the static characteristics of the nucleons. When considering electromagnetic characteristics we shall use the matrix element (13) and write it down in two-component spinors:

$$(\omega_{S'}^{N})^{+}\left[F_1(\vec{K}_\perp^2)+iF_2(\vec{K}_\perp^2)K_{\perp a}\varepsilon_{ab}\sigma_b\right]\omega_S^N = \tag{52}$$

$$=3\sum_{\substack{s_a,s_b\\ s_c,s_c'}}\int d\Gamma\left[\psi^{S'}_{s_a,s_b,s_c'}(\xi,\eta,\vec{q}_\perp,\vec{Q}_\perp+\eta\vec{K}_\perp)\right]^{+}\cdot$$

$$\cdot\,\omega^{+}_{s_c'}\left(Q_c+i\frac{\varkappa_c}{2m_c}\vec{K}_{\perp a}\varepsilon_{ab}\sigma_b\right)\omega_{s_c}\Psi^{S}_{s_a,s_b,s_c}(\xi,\eta,\vec{q}_\perp,\vec{Q}_\perp).$$

In (52) $F_1(\vec{K}_\perp^2)$ and $F_2(\vec{K}_\perp^2)$ are nucleon Pauli formfactors, the vertex functions Ψ are determined by the relations (16, 17, 50). The relation (52) is equivalent to the analogous one obtained in [8] in the light-front dynamics.

If one expand (52) in powers of $K_\perp$, then the terms of this expansion containing $K_\perp$ in the zero, first, second and third powers determine respectively the nucleon charge $F_1(0)$, the nucleon anomalous magnetic moment $F_2(0)$ as well as the radii corresponding to the formactors $F_1(\vec{K}_\perp^2)$ and $F_2(\vec{K}_\perp^2)$. The quantities $F_2(0)$ have been calculated in Refs [8, 10] and have the following form:

$$\frac{\varkappa_p}{2M_p}=\frac{2X^c+Z(4\varkappa_u-\varkappa_d)}{6m_q}\,;\quad \frac{\varkappa_n}{2M_n}=\frac{-\frac{2}{3}(2X^c+X^a)+Z(4\varkappa_d-\varkappa_u)}{6m_q}\,, \tag{53}$$

where

$$X^c = \int \frac{\gamma}{2} \frac{3m_q}{M_0} \left|\phi(M_0^2)\right|^2 d\Gamma, \quad X^a = \int \frac{\gamma(1-\eta)}{\eta} \frac{3m_q}{M_0} \left|\phi(M_0^2)\right|^2 d\Gamma, \quad (54)$$

$$Z = \int \frac{\beta^2}{\vec{Q}_\perp^2 + \beta^2} \left|\phi(M_0^2)\right|^2 d\Gamma, \qquad \gamma = 2\,\frac{\eta M_0 \beta - \frac{\vec{Q}_\perp^2}{2}}{Q_\perp^2 + \beta^2},$$

$$\beta = m_q + (1-\eta) M_0 .$$

The calculations and final formulae for the electromagnetic radii are too cumbersome and it seems impossible to present them in this work. To illustrate the role of the relativistic effects we shall give the first two terms of the expansion of our results in powers of $\left(\frac{\alpha}{m_q}\right)^2$. Note that parameter α characterizes the quark mean square momenta in nucleons: $\langle Q^2 \rangle = \frac{4}{3} \langle q^2 \rangle = \alpha^2$. In the nonrelativistic limit $\alpha \to 0$

These expansions have the form:

$$F_2^p(0) = \frac{1}{3m_q}\left(1 - \frac{4}{12}\frac{\alpha^2}{m_q^2} + \dots\right), \quad F_2^n(0) = -\frac{1}{3m_q}\left(1 - \frac{5}{12}\frac{\alpha^2}{m_q^2} + \dots\right) \quad (55)$$

$$R^2_{F_1 p} = \frac{1}{\alpha^2} + \frac{10}{12 m_q^2} + \dots, \qquad R^2_{F_1 n} = \frac{1}{3m_q^2} + \dots,$$

$$R^2_{F_2 p} = \frac{1}{\alpha^2} + \frac{15}{12 m_q^2} + \dots, \qquad R^2_{F_2 n} = \frac{1}{\alpha^2} + \frac{16}{12 m_q^2} + \dots .$$

For simplicity we have not given the contribution of the quark anomalous magnetic moments. From the expansions obtained one can see that the relativistic corrections do not violate essentially the nonrelativistic quark model predictions for the nucleon magnetic moment ratio and for the relations between nucleon electromagnetic radii.

The ratio G_A/G_V for nucleon one can easily derive replacing the electromagnetic current in (13) by charged axial current:

$$G_A/G_V = \frac{5}{3}\cdot T, \tag{56}$$

where

$$T = \int \frac{\beta^2 - Q_\perp^2}{\beta^2 + Q_\perp^2}\left|\phi(M_\circ^2)\right|^2 d\Gamma \tag{57}$$

In the nonrelativistic limit the well-known quark model prediction $\frac{G_A}{G_V} = \frac{5}{3}$, which exceeds the experimental value, follows from (56). The relativistic effects change this value into the needed side decreasing the quark model prediction, the first two terms of the formulae (56) expansion in powers of $\left(\frac{\alpha}{m_q}\right)^2$ being of the form:

$$\frac{G_A}{G_V} = \frac{5}{3}\left(1 - \frac{\alpha^2}{3m_q^2} + \ldots\right). \tag{58}$$

The quantitative analysis of the experimental data by our entire formulae (without expanding them in $\left(\frac{\alpha}{m_q}\right)^2$) we have done using the linear square fit. The fitting parameters α, m_q, $\varkappa$ and $\tilde{\varkappa}$, where $\varkappa$ and $\tilde{\varkappa}$ are determined by the relation $\varkappa_c = Q_c\varkappa + \tilde{\varkappa}$ turned out equal to

$$\alpha = 379\pm61\ \text{MeV};\quad m_q = 271\pm28\ \text{MeV};\quad \varkappa = 0.071\pm0.040,\quad \tilde{\varkappa} = -0.035\pm0.015 \tag{59}$$

Note that for the magnetic moments and the ratio $\frac{G_A}{G_V}$ which are measured with high precision we took 3% errors instead of experimental ones. Thus we have taken into account the uncertainty of our results due to the assumptions made, in particular, due to the assumptions on the form of the functions $\varphi(p_a, p_b, p_c)$ and $\phi(M_\circ^2)$, and also due to the assumption that nucleon is a pure $[56, 0^+]$ state without admixture of other multiplets.

The value of $\widetilde{\mathcal{H}}$ (59) agrees with its value obtained from the previous data on the ratio of $\rho \to \pi\gamma$ and $\omega \to \pi\gamma$ decay widths [10, 16]. The new data on $\rho \to \pi\gamma$ [17] indicate to a smaller value of $|\widetilde{\mathcal{H}}|$, but do not exclude the value obtained by us. It is interesting that our parameters α and m_q (59) agree rather well with their values obtained from the mass formulae [5].

The results of our analysis are presented in Table 1. The radii given in this Table correspond to the Sachs formfactors. To demonstrate the role of relativistic effects and quark a.m.m. we give separately the following contributions: (a) - contribution corresponding to the nonrelativistic approximation in eqs.(55, 58), (b) - total contribution of all relativistic corrections in (55, 58), (c) - contribution of quark a.m.m. From the Table (see also eq.(55)) one can see that the relativistic effects are not small, but they nowhere violate the relations predicted by the nonrelativistic quark model. For $\frac{G_A}{G_V}$ the relativistic effects lead to the agreement of the theory with the experiment. Emphasize that we received the value of the neutron electric radius $R^2_{G^n_E}$ assuming that nucleon is pure $[56, 0^+]$ state, whereas in other works the nonzero value of $R^2_{G^n_E}$ had been obtained only by taking into account the admixture of other states (see. e.g. [18]).

Thus the treatment of the nucleon as bound state of the relativistic quarks which is in the basis of our model allows one to describe self-consistently all static characteristics of the nucleon. The success of the nonrelativistic model when describing the ratios $\frac{\mu_p}{\mu_n}$, $\frac{G_A}{G_V}$ and relations between electromagnetic radii of nucleons is apparently of casual nature and is due to the fact that for these ratios (as distinct from the quantities μ_p, μ_n... themselves) relativistic effects are numerically small.

5. Amplitudes of $N^* \to N\gamma$ and $N^* \to N\pi$ decays.
Selection rules for matrix elements.

Consider linear in $K_\perp$ terms of the matrix elements (13') expansion in powers of $K_\perp$. Making use of the transition current $\langle N | J_\mu | N^* \rangle$ explicit form (e.g., from Ref.[21]) one can show that these terms can be written in the form:

$$\frac{1}{2P} \langle N, \lambda - 1 | J_{0,3} | N^*, \lambda \rangle \Big|_{P \to \infty} \equiv M^{\lambda}_{0,3} \frac{K_1 - iK_2}{\sqrt{2}} , \tag{60}$$

where λ and $\lambda - 1$ are respectively helicities of resonance and nucleon. The amplitudes $M^{\lambda}_{0,3}$ defined in such a way are connected with the formfactors G_1 and G_2 introduced in [21] by the relations:

$$J^P = \frac{1}{2}^+ : \qquad M^{\frac{1}{2}}_{0,3} = G_2 \frac{M_N - M^*}{\sqrt{2}} , \tag{61}$$

$$J^P = \frac{3}{2}^+ : \qquad M^{\frac{1}{2}}_{0,3} = -\frac{1}{\sqrt{3}} \left(-\frac{M_N}{M^*} G_1 + \frac{M^* - M_N}{2} G_2 \right) ,$$

$$M^{\frac{3}{2}}_{0,3} = -G_1 + \frac{M_N - M^*}{2} G_2 ,$$

where M^*, J and P - are mass, spin and parity of the resonance. For $P = -1$ the replacement $M^* \to -M^*$ must be done in (61). For the decay widths of $N^* \to N\gamma$ we obtain

$$\Gamma(N^*_J \to N\gamma) = \frac{4\alpha}{2J+1} q_c^3 \sum_{\lambda = \frac{1}{2}, \frac{3}{2}} \left| M^{\lambda}_{0,3} \right|^2 , \tag{62}$$

$\alpha = \frac{1}{137}$, q_c is the photon momentum in the resonance rest frame. The amplitudes $M^{\lambda}_{0,3}$ are connected with the conventional ones $A^{p,n}_{\lambda}$, for which the experimental data are usually given, by the relation:

$$A_\lambda^{p,n} = \sqrt{\alpha\pi \frac{M^{*2} - M_N^2}{M_N}}\; M_{0,3}^\lambda \,\mathrm{sign}\, A_\pi \left(N^{*+,0} \rightarrow n\pi^{+,0}\right). \tag{63}$$

This relation includes the sign of the amplitude of $N^* \rightarrow N\pi$ since quantities $A_\lambda^{p,n}$ are determined from the experimental data on the reaction $\gamma N \rightarrow N\pi$ and the sign of $N^* \rightarrow N\pi$ transition amplitude enters the $A_\lambda^{p,n}$ definition.

In our approach in accordance with (13') the amplitudes $M_{0,3}^\lambda$ are determined by the relation:

$$M_{0,3}^\lambda \frac{K_1 - iK_2}{\sqrt{2}} = 3\sum_{\substack{S_a,S_b,\\ S_c,S_c'}} \int d\Gamma \left[\Psi^{\lambda-1}_{N,S_a,S_b,S_c'}(\xi,\eta,\vec{q}_\perp,\vec{Q}_\perp)\right]^+ \tag{64}$$

$$\omega^+_{S_c'}\left(Q_c\eta K_{\perp\ell}\frac{\overleftarrow{\partial}}{\partial Q_{\perp\ell}} + i\frac{\varkappa_c}{2m_c}K_{\perp\ell}\varepsilon_{\ell k}\sigma_k^c\right)\omega_{S_c}\Psi^\lambda_{N^*,S_a,S_b,S_c}(\xi,\eta,\vec{q}_\perp,\vec{Q}_\perp),$$

where the derivative $\frac{\partial}{\partial Q_{\perp\ell}}$ acts on the nucleon vertex function.

Making use of PCAC one can express the amplitude $N^* \rightarrow N\pi$ in terms of the charged axial current matrix element in the $P \rightarrow \infty$ frame:

$$A_\pi(N^{*+} \rightarrow n\pi^+) = \frac{1}{2\sqrt{2}P}\left\langle n,\lambda=\tfrac{1}{2}\right|\left(J^{axial}\right)^{1-i2}_{0,3}\left|N^{*+},\lambda=\tfrac{1}{2}\right\rangle\Big|_{P\rightarrow\infty}, \tag{65}$$

the $N^{*+} \rightarrow n\pi^+$ decay width being of the form

$$\Gamma(N^{*+} \rightarrow n\pi^+) = \frac{1}{2\pi F_\pi^2}\;\frac{q_{r_c}^\pi}{2J+1}\;\frac{(M^{*2} - M_N^2)^2}{M^{*2}}\left|A_\pi(N^{*+} \rightarrow n\pi^+)\right|^2, \tag{66}$$

where $q_{r_c}^\pi$ is the pion momentum in the resonance rest frame, F_π = 135 MeV is the $\pi \rightarrow \mu\nu$ decay constant. In our approach in accordance with (65) and (13') we have

$$A_{\pi}(N^{*+}\to n\pi^{+}) = \frac{3}{\sqrt{2}}\sum_{s_a,s_b,s_c}\int d\Gamma\left[\Psi^{\lambda=\frac{1}{2}}_{n,s_a,s_b,s_c}(\xi,\eta,\vec{q}_{\perp},\vec{Q}_{\perp})\right]^{+}\cdot \tag{67}$$

$$\sigma_3^c\,\tau_-^c\,\Psi^{\lambda=\frac{1}{2}}_{N^{*+},s_a,s_b,s_c},$$

where τ_- is operator transforming $u\to d$.

Consider selection rules for the matrix elements (64). In our approach the functions Ψ in (64) written in terms of spinors W_S (32) have the well known SU(6)xO(3) structure. Let us write them schematically in the form:

$$\left[\Psi_{N,N^*}\right]^{+} = \mathcal{U}(K_a)\,\mathcal{U}(K_b)\,\mathcal{U}(K_c)\left[\Psi^{SU(6)\times O(3)}_{N,N^*}\right]^{+}, \tag{68}$$

where $\mathcal{U}(K_a)$, $\mathcal{U}(K_b)$, $\mathcal{U}(K_c)$ are the Melosh matrices determined according to (23). When the differentiation operator $\frac{\overleftarrow{\partial}}{\partial Q_{\perp\ell}}$ acts on the matrix $\mathcal{U}(K_c)$ the Melosh matrices for a- and b- quarks are contracted and the selection rules are determined by the operator:

$$\mathcal{U}(K_c)\frac{\overleftarrow{\partial}}{\partial Q_{\perp\ell}}\left[\mathcal{U}(K_c)\right]^{+} = \frac{2Q_\ell + i\sigma_k\varepsilon_{k\ell}[m_q+M_o(1-\eta)] - i\sigma_3 Q_k\varepsilon_{k\ell} - \frac{i(\vec{\sigma}\vec{\varepsilon}\vec{Q}_{\perp})Q_{\perp\ell}}{M_o\eta}}{\vec{Q}_{\perp}^2+[m_q+M_o(1-\eta)]^2}. \tag{69}$$

Keeping in mind that the operator Q_c acts on the same quark c we obtain the following selection rules corresponding in succession to the terms of eq.(69)

$$\{35\}\begin{cases}\Delta S=0,\ \Delta S_z=0,\ \Delta L_z=\pm 1, & (70.1)\\ \Delta S=1,\ \Delta S_z=\pm 1,\ \Delta L_z=0, & (70.2)\\ \Delta S=1,\ \Delta S_z=0,\ \Delta L_z=\pm 1, & (70.3)\\ \Delta S=1,\ \Delta S_z=\pm 1,\ \Delta L_z=\pm 2, & (70.4)\end{cases}$$

where S and L are total spin and orbital angular momentum of quarks in c.m.s. Note that the selection rules (70.2) and (70.3) hold also for that part of quark a.m.m. operator which is proportional to the quark charge operator.

The selection rules (70) coincide with those of [22] following from the transformation between generators of the $SU(6)_{w,current}$ and $SU(6)_{w,constituent}$ groups obtained by Melosh in the free quark model. However in our approach the matrix element (64) contains the additional to (70) terms which arise when the operator $\frac{\overleftarrow{\partial}}{\partial Q_{\perp\ell}}$ acts on the matrices $\mathcal{U}(K_a)$ and $\mathcal{U}(K_b)$. Emphasize that these terms do not belong to any representation of SU(6) group, since in these terms the charge operator Q_c and the differentiation operator $\frac{\overleftarrow{\partial}}{\partial Q_{\perp\ell}}$ act on different quarks.

Consider the role of these additional to the selection rules (70) terms on the example of the nucleon magnetic moments. The relation (64) for nucleons determines the nucleon anomalous magnetic moments; the calculations results are given in (54). The quantities X^a and X^c correspond to acting of the differentiation operator $\frac{\partial}{\partial Q_{\perp\ell}}$ on the Melosh matrices $\mathcal{U}(K_a)$ and $\mathcal{U}(K_c)$, respectively. Note that in the nonrelativistic limit $X^a = X^c = 1$. The result of the acting of $\frac{\partial}{\partial Q_{\perp\ell}}$ on $\mathcal{U}(K_c)$ brings to the selection rule (70.2) and to the corresponding to it value of the ratio $\frac{\varkappa_p}{\varkappa_n} = -\frac{3}{2}$ which contradicts the experiment. The agreement with the experiment can be obtained only by taking into account X^a, i.e. terms which cannot be obtained in the Melosh approach. Note that the value of the ratio $\frac{\varkappa_p}{\varkappa_n} = -\frac{3}{2}$ following from the selection rule (70.2) which takes place for the nucleon a.m.m. was misreferred by Melosh to the nucleon total magnetic moments [13]. The doubtfulness of this result of Melosh was discussed in Ref. [23].

The selection rules obtained in our approach for the matrix elements (67) coincide with the analogous ones[22] following from the Melosh transformation between current and constituent quarks. This is due to the fact that in the matrix elements (67) there are no operators which act simultaneously on the quarks (a, b) and c. These selection rules are as follows:

$$\Delta S = 1, \quad \Delta S_z = 0, \quad \Delta L_z = 0, \tag{71.1}$$

$$\Delta S = 1, \quad \Delta S_z = \pm 1, \quad \Delta L_z = \mp 1. \tag{71.2}$$

6. $P_{33}\,(1232) \to N\gamma$ decay

The amplitudes of this decay defined by the relation (64) and calculated analogously to a.m.m. of nucleon are equal to

$$M_{0,3}^{\frac{1}{2}} = -\frac{\varkappa_n}{2M_n} + z\,\frac{\tilde{\varkappa}}{2m_q}\,, \tag{72}$$

$$M_{0,3}^{\frac{3}{2}} = \sqrt{3}\ M_{0,3}^{\frac{1}{2}}\,.$$

From (72) directly follows the suppression of the amplitude E_{1+}:

$$E_{1+} = 0\,, \tag{73}$$

which coincides with the well-known result of the nonrelativistic quark model and agrees well with the experiment. Thus relativistic effects do not violate this result.

To obtain the signs of the amplitudes $A_\lambda^{p,n}$ we have used the relations (63, 65, 67). So we have

$$\mathrm{sign}\, A_\pi\,(\Delta^+ \to n\pi^+) = -1\,. \tag{74}$$

The numerical values of the amplitudes A_λ^p following from (63), (64) and (74) corresponding to the parameters (59) are given in Table 2. For comparison we have presented in the Table the predictions of the nonrelativistic quark model which differ considerably from experimental data. This discrepancy is usually treated as a serious difficulty of nonrelativistic quark model (see, e.g. [24]). Our amplitudes $A^p_{\frac{1}{2}}$ and $A^p_{\frac{3}{2}}$ are essentially closer to the experimental values.

7. $P_{11}(1470) \to N\gamma$ and $P_{11}(1780) \to N\gamma$ decays

We have calculated the amplitudes of these decays assuming that $P_{11}(1470)$ and $P_{11}(1780)$ resonances are the members of $[56, 0^+]_{r}$ and $[70, 0^+]$ multiplets, respectively. The vertex function of the $P_{11}(1470)$ resonance we have written down proceeding from the analogy with the nonrelativistic oscillator model:

$$\Psi_{P_{11}(1470)} = N\left(M_0^2 - \beta\right)\Psi_N \quad , \tag{75}$$

where parameter β can be found out from the orthogonality condition of the functions (17) and (75), N is the normalization parameter.

The obtained by us amplitudes $M^{\frac{1}{2}}_{0,3}$ for $P_{11}(1470) \to N\gamma$ decays are as follows:

$$M^{\frac{1}{2}}_{0,3}(p) = -\sqrt{2}\,\frac{2X^c_R + Z_R(4\varkappa_u - \varkappa_d)}{6m_q} \quad , \tag{76}$$

$$M^{\frac{1}{2}}_{0,3}(n) = -\sqrt{2}\,\frac{-\frac{2}{3}(2X^c_R + X^a_R) + Z_R(4\varkappa_d - \varkappa_u)}{6m_q}$$

For the amplitude $A_\pi\left(P^+_{11}(1470) \to n\pi^+\right)$ we obtain

$$A_\pi\left(P^+_{11}(1470) \to n\pi^+\right) = \frac{5}{3\sqrt{2}}\,T_R \,. \tag{77}$$

The quantities X^a_R, X^c_R, Z_R and T_R differ from the corresponding quantities for the nucleon (54, 57) by factor $N(M_\circ^2-\beta)$ in the integrands. The final formulae for $P_{11}(1780)\rightarrow N\gamma$ decay amplitudes are cumbersome, so we give for them only numerical results.

The numerical results for $A_\lambda^{p,n}$ amplitudes at the values of parameters (59) are given in Table 2. Note that as contrary to the predictions of nonrelativistic quark model, the obtained by us signs of the amplitudes $A_{\frac{1}{2}}^p$ and $A_{\frac{1}{2}}^n$ for $P_{11}(1470)\rightarrow N\gamma$ decay agree with the experiment. It is known that the discrepancy of these amplitude signs with the experiment is considered as a serious difficulty of the nonrelativistic quark model (see, e.g. [24]). In our approach we have obtained correct signs, however the numerical values of our amplitudes are smaller than the experimental ones. This apparently indicates to the mixing of this resonance with the nucleon. Note that here in order to obtain the experimental values of the amplitudes $A_\lambda^{p,n}(P_{11}(1470)\rightarrow N\gamma)$ it is enough to introduce the small mixing angle with the nucleon ($\cong 5^0$).

The amplitudes $A_{\frac{1}{2}}^{p,n}$ obtained by us for the transition $P_{11}(1470)\rightarrow N\gamma$ satisfy the relation $A_{\frac{1}{2}}^p \cong -A_{\frac{1}{2}}^n$ which does not coincide with the prediction $A_{\frac{1}{2}}^p / A_{\frac{1}{2}}^n = -\frac{3}{2}$ [23] obtained in the Melosh approach. This is due to the fact that, as is shown in Sec.5, the Melosh approach corresponds to the contribution in (76) of X_R^c only.

The experimental data on $P_{11}(1780)\rightarrow N\gamma$ transition are contradictory, therefore we cannot make any definite conclusions on the obtained by us results for this transition.

8. $[70,1^-]\rightarrow N\gamma$ decays

For the amplitudes $M_{0,3}^\lambda$ of these decays we have found from (64):

$$M_{0,3}^\lambda = \sum_{i=1}^{9} a_i(N^*,\lambda)A_i\ , \qquad (78)$$

where the factors $\alpha_i(N^*, \lambda)$ are given in Table 3, and quantities A_i are equal to

$$A_i = \frac{1}{\sqrt{N_N N_R}} \int d\Gamma' B_i \Phi(M_0^2) \Phi_{N*}(M_0^2), \tag{79}$$

$$d\Gamma' = \frac{M_{ab}}{\varepsilon_a \varepsilon_b} \frac{M_0}{E_{ab} E_c} q^2 dq\, Q^2 dQ\, d\cos\theta_Q, \tag{80}$$

$$N_N = \int |\Phi(M_0^2)|^2 d\Gamma', \quad N_R = \int |\Phi_{N*}(M_0^2)|^2 d\Gamma',$$

where $\Phi_{N*}(M_0^2)$ are the radial part of the N^* resonance vertex functions. The integrands B_i have the following form

$$B_1 = -\frac{Q_\perp^2}{\sqrt{2}\alpha^2} \frac{M_0}{E_c - Q_z}, \tag{81}$$

$$B_2 = \left\{ 3\sqrt{2} \frac{Q_z}{M_0} \left[\frac{Q_\perp^2}{2} - R(E_{ab} + Q_z) \right] + \frac{3}{2} \frac{Q_z R^2}{m_q} \varkappa \right\} / R_1,$$

$$B_3 = -\frac{3}{\sqrt{2}} \frac{E_{ab} + Q_z}{M_0} \frac{Q_\perp^2}{R_1} + \frac{3}{2} \frac{R \cdot Q_\perp^2}{R_1 \cdot m_q} \varkappa,$$

$$B_4 = -\frac{3}{2\sqrt{2}} \frac{E_c - Q_z}{M_0} \frac{Q_\perp^2}{R_1}, \quad B_5 = 3B_4,$$

$$B_6 = 3\sqrt{2}(C_1 + C_2), \quad B_7 = (-C_1 + 3C_2)3\sqrt{2},$$

$$B_8 = \frac{Q_\perp^2 Q_z}{R_1 m_q} \tilde{\varkappa}, \quad B_9 = \frac{R \cdot Q_\perp^2}{R_1 m_q} \tilde{\varkappa},$$

where

$$C_1 = \frac{\mathcal{E}_a}{4q} Q_\perp^2 \left(2\mathcal{E}_a q - m_q^2 \ln \frac{\mathcal{E}_a + q}{\mathcal{E}_a - q}\right) \frac{E_c - Q_z}{M_0 (E_{ab} + Q_z)^2 R_1} , \qquad (81')$$

$$C_2 = - \frac{(E_c - Q_z)\, Q_z}{2M_0 (E_{ab} + Q_z) R_1} \left[R(E_{ab} + Q_z) - \frac{Q_\perp^2}{2}\right],$$

$$R = m_q + E_c - Q_z , \qquad R_1 = R^2 + Q_\perp^2 .$$

Note that the first three terms in (78) correspond to the selection rules (70.1), (70.2) and (70.3), respectively, and factors a_1, a_2, a_3 coincide with the corresponding ones of Refs.[23, 25] obtained in the Melosh approach. The remaining terms in (78) correspond to the action of the operator $\frac{\partial}{\partial Q_{\perp\ell}}$ on the matrices $\mathcal{U}(K_{a,b})$ in (64) as well as to the contribution of the quarks a.m.m. not proportional to their charge. As was shown in Sec.5, the analogous terms cannot be obtained in the Melosh approach.

For the matrix element of axial current (67) we find:

$$A_\pi (N^{*+} \to n\pi^+) = a(N^*)\, A + b(N^*)\, B , \qquad (82)$$

where factors a and b are given in Table 4, and quantities A and B are equal to

$$A = - \frac{6\sqrt{2}}{\sqrt{N_N N_R}} \int d\Gamma' \frac{Q_\perp^2\, Q_z}{R_1} \phi(M_0^2)\, \phi_{N*}(M_0^2), \qquad (83)$$

$$B = - \frac{3\sqrt{2}}{\sqrt{N_N N_R}} \int d\Gamma' \frac{R\, Q_\perp^2}{R_1} \phi(M_0^2)\, \phi_{N*}(M_0^2).$$

The first and second terms correspond to the selection rules (71.1) and (71.2), respectively, the factors a and b coinciding with the corresponding ones of Ref.[26] obtained in the Melosh approach.

When getting numerical results we assume the analogous to (17) form of the radial part of the N^* resonance vertex function. The resonances under question being members of the multiplet $[70,1^-]$, the parameter α_{N^*} for these resonances differs, generally speaking, from the corresponding nucleon parameter α (58). To find α_{N^*} we use the well established data on $A^p_{\frac{3}{2}}$ amplitude of the $D_{13}(1520)\to N\gamma$ transition. The value of this amplitude can be obtained for two values of the parameter

$$\alpha^{(1)}_{N^*} = 228 \text{ MeV}, \qquad \alpha^{(2)}_{N^*} = 813 \text{ MeV}. \tag{84}$$

As to the other amplitudes of $[70,1^-]\to N\gamma$ transitions, we predict them. The results are almost the same for $\alpha^{(1)}_{N^*}$ and $\alpha^{(2)}_{N^*}$.

In Table 2 we present the results corresponding to the parameters (58) in the $\alpha^{(1)}_{N^*}$ case. One can see from the Table that the obtained results agree well with the experiment. Let us emphasize the following features of our results: a) we predict the signs of the amplitudes $A^{p,n}_{\lambda}$ of $[70,1^-]\to N\gamma$ decays which agree well with the experiment, b) our amplitudes of $S_{11}(1535)\to N\gamma$ and $D_{13}(1520)\to N\gamma$ decays obey the relation $A^p_{\frac{1}{2}}(S_{11}) < A^p_{\frac{3}{2}}(D_{13})$ which agrees with the experiment.

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Table 1

	μ_p nuclear magneton	μ_n nuclear magneton	$\frac{G_A}{G_V}$	$R^2_{G_E^p}$ [F^2]	$R^2_{G_M^p}$ [F^2]	$R^2_{G_E^n}$ [F^2]	$R^2_{G_M^n}$ [F^2]
Our results	2.811	-1.848	1.222	0.657	0.681	-0.094	0.610
(a)	3.308	-2.308	1.667	0.275	0.326	0	0.331
(b)	-0.604	0.706	-0.445	0.408	0.316	-0.049	0.271
(c)	0.107	-0.246		-0.026	0.039	-0.045	0.008
Exsperi-ment	2.793	-1.913	1.254 ±0.007	0.667 ±0.013	0.689 ±0.033	-0.121 ±0.004	0.689 ±0.049
	[19]	[19]	[19]	[15]	[15]	[20]	[15]

Table 2

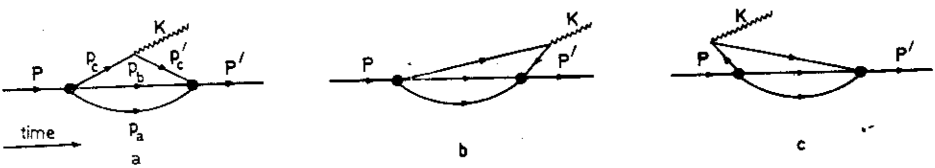
Multiplet	Resonance	Amplitude	Experiment C [27]	Experiment BCP [28]	Experiment Tokyo [29]	Quark models predictions FKR [30]	Quark models predictions KO [31]	Quark models predictions KI [32]	Our pre-dic-tions
$[56,0^+]$	$P_{33}(1232)$	$A^p_{1/2}$	-130	-142	-145	-108	-101	-103	-121
		$A^p_{3/2}$	-248	-271	-261	-187	-176	-179	-210
$[56,0^+]_2$	$P_{11}(1470)$	$A^p_{1/2}$	-70	-75	-66	27	-5	-24	-28
		$A^n_{1/2}$	44	59	19	-18	4	16	33
$[70,1^-]$	$S_{11}(1535)$	$A^p_{1/2}$	32	82	80	156	97	147	98
		$A^n_{1/2}$	-88	-112	-75	-108	-101	-119	-89
	$S_{11}(1700)$	$A^p_{1/2}$	44	48	61	0	0	88	4
		$A^n_{1/2}$	-103	-45	8	30	4	-35	9
	$D_{13}(1520)$	$A^p_{1/2}$	-9	-16	-32	-34	6	-23	17
		$A^p_{3/2}$	162	157	162	109	174	128	157
		$A^n_{1/2}$	-67	-55	-71	-31	-52	-45	-50
		$A^n_{3/2}$	-133	-141	-148	-109	-144	-122	-142
	$D_{13}(1700)$	$A^p_{1/2}$	-12	-33	-29	0	0	-7	-19
		$A^p_{3/2}$	-12	-14	14	0	0	11	2
		$A^n_{1/2}$	81	50	-55	-10	15	-15	13
		$A^n_{3/2}$	107	35	-35	-40	-17	-76	-30
	$S_{31}(1650)$	$A^p_{1/2}$	44	34	-26	47	86	59	81
	$D_{33}(1670)$	$A^p_{1/2}$	101	130	130	88	90	100	117
		$A^p_{3/2}$	116	98	50	84	91	105	151
$[70,0^+]$	$P_{11}(1780)$	$A^p_{1/2}$	62	1	-12	-40	-7	-47	-30
		$A^n_{1/2}$	-20	-28	11	10	2	-21	45

Table 3

transition	a_1	a_2	a_3	a_4	a_5	a_6	a_7	a_8	a_9	helicity
$S_{11}(1535)\to\gamma p$	$-1/6$	$-1/6$	$1/6$	$1/9$	0	0	$-1/9$	$1/2$	$1/2$	$1/2$
$\to\gamma n$	$1/6$	$1/18$	$-1/18$	0	$-1/9$	$1/9$	$1/18$	$1/2$	$1/2$	$1/2$
$S_{11}(1700)\to\gamma p$	0	0	0	$-1/9$	0	0	$-1/9$	$1/4$	$1/4$	$1/2$
$\to\gamma n$	0	$-1/18$	$1/18$	$1/18$	$-1/18$	$1/18$	$1/18$	$1/4$	$1/4$	$1/2$
$D_{13}(1520)\to\gamma p$	$-\sqrt{2}/12$	$\sqrt{2}/6$	$\sqrt{2}/12$	$\sqrt{2}/18$	0	0	$\sqrt{2}/9$	$-\sqrt{2}/2$	$\sqrt{2}/4$	$1/2$
	$-\sqrt{6}/12$	0	$-\sqrt{6}/12$	$-\sqrt{6}/18$	0	0	0	0	$-\sqrt{6}/4$	$3/2$
$\to\gamma n$	$\sqrt{2}/12$	$-\sqrt{2}/18$	$-\sqrt{2}/36$	0	$-\sqrt{2}/18$	$-\sqrt{2}/9$	0	$-\sqrt{2}/2$	$\sqrt{2}/4$	$1/2$
	$\sqrt{6}/12$	0	$\sqrt{6}/36$	0	$\sqrt{6}/18$	0	0	0	$-\sqrt{6}/4$	$3/2$
$D_{13}(1700)\to\gamma p$	0	0	0	$-2\sqrt{5}/45$	0	0	$-\sqrt{5}/45$	$\sqrt{5}/20$	$\sqrt{5}/5$	$1/2$
	0	0	0	$\sqrt{15}/45$	0	0	$-\sqrt{15}/15$	$3\sqrt{15}/20$	$\sqrt{15}/10$	$3/2$
$\to\gamma n$	0	$-\sqrt{5}/90$	$2\sqrt{5}/45$	$-4\sqrt{5}/45$	$-2\sqrt{3}/27$	$\sqrt{5}/90$	$-\sqrt{5}/90$	$\sqrt{5}/20$	$\sqrt{5}/5$	$1/2$
	0	$-\sqrt{15}/30$	$\sqrt{15}/45$	$-2\sqrt{15}/45$	$-\sqrt{15}/45$	$\sqrt{15}/30$	$\sqrt{15}/30$	$\sqrt{15}/100$	$\sqrt{15}/10$	$3/2$
$S_{31}(1650)\to\gamma p$	$-1/6$	$1/18$	$-1/18$	$1/9$	$2/9$	$-2/9$	$-1/9$	0	0	$1/2$
$D_{33}(1670)\to\gamma p$	$-\sqrt{2}/12$	$-\sqrt{2}/18$	$-\sqrt{2}/36$	$\sqrt{2}/18$	$\sqrt{2}/9$	$2\sqrt{2}/9$	$\sqrt{2}/9$	0	0	$1/2$
	$-\sqrt{6}/12$	0	$\sqrt{6}/36$	$-\sqrt{6}/18$	$-\sqrt{6}/9$	0	0	0	0	$3/2$

Table 4

Resonance	$S_{11}(1535)$	$S_{11}(1700)$	$D_{13}(1520)$	$D_{13}(1700)$	$S_{31}(1650)$	$D_{33}(1670)$
	$^2[8]_{1/2}$	$^4[8]_{1/2}$	$^2[8]_{3/2}$	$^4[8]_{3/2}$	$^2[10]_{1/2}$	$^2[10]_{3/2}$
α	$-\frac{1}{9}$	$-\frac{1}{18}$	$\frac{\sqrt{2}}{9}$	$-\frac{\sqrt{5}}{90}$	$-\frac{1}{36}$	$\frac{\sqrt{2}}{36}$
β	$-\frac{2}{9}$	$-\frac{1}{9}$	$-\frac{\sqrt{2}}{9}$	$\frac{\sqrt{5}}{90}$	$-\frac{1}{18}$	$-\frac{\sqrt{2}}{36}$



Figure

REFERENCES

1. Lipkin H.J. Phys.Rep., 1973, 3C, 173.
2. Shekhter V.M. Proc. of the IV Winter School of LNPI on Nuclear Physics and Elementary Particles, 1976, v.1, p.38.
3. Anisovich V.V., Shabelsky Yu.M., Shekhter V.M. Yad.Fiz. 1978, 28, 1063.
4. Von Gehlen G. Proc. of the 6-th Intern. Symp. on Electron and Photon Interactions at High Energies, Bonn, 1973, p.117.
5. Isgur N., Karl G. Phys.Rev. 1978, D18, 4187; Phys.Rev. 1979, D19, 2653, Phys.Rev. 1979, D20, 1191.
6. Terent'ev M.V. Yad.Fiz. 1976, 24, 207; ITEP-preprint 1976, N.6.
7. Berestetski V.B., Terent'ev M.V. Yad.Fiz. 1976, 24, 1044.
8. Berestetski V.B., Terent'ev M.V. Yad.Fiz. 1977, 25, 653.
9. Kondratyuk L.A., Terent'ev M.V. Yad.Fiz. 1980, 31, 1087.
10. Aznauryan I.G., Ter-Isaakyan N.L. Yad.Fiz. 1980, 31, 1680.
11. Leutwyler H., Stern J. Ann. of Phys. 1979, 112, 107.
12. Frankfurt L., Strickman M. Nucl.Phys. 1979, B148, 107.
13. Melosh H.J. Phys.Rev. 1974, D16, 1095.
14. Cashmore R.J. Proc. of the 19-th Intern. Conf. on High Energy Physics, Tokyo, 1978, p.811.
15. Kravtsov A.V., Nemenov L.L. Yad.Fiz. 1977, 26, 145.
16. Kamal A.N. Phys.Rev. 1978, D18, 3512.
17. Berg D., Chandlee C., Cihangir S. et al. Phys.Rev.Lett. 1980, 44, 706.
18. Isgur N., Karl G., Koniuk R. Phys.Rev.Lett. 1978, 41, 1269.
19. Review of Particles Properties. Rev. of Modern Phys. 1980, 52.
20. Koestler L., Nistler W., Waschkowski W. Phys.Rev.Lett. 1976, 36, 1021.
21. Devenish R.C.E., Eisenschitz T.S., Korner J.G. Preprint DESY 75/48, 1975.

22. Babcock J., Rosner J.L. Ann. of Phys. 1976, 96, 191; Hey A.J.G., Weyers J. Phys.Lett., 1974, 48B, 69.

23. Gilman F.J., Karliner I. Phys.Rev. 1974, D10, 2194.

24. Kajikawa R. Nagoya preprint DPNU 31-80, September 1980.

25. Aznauryan I.G., Bagdasaryan A.S. Izv. Akad. Nauk Arm.SSR, Fizika, 1977, 12, 416.

26. Gilman F.J., Kugler M., Meshkov S. Phys.Rev. 1974, D9, 715.

27. Crawford R.L. Nucl.Phys. 1975, B97, 125.

28. Barbour I.M., Crawford R.L., Parsons N.H. Nucl.Phys. 1978, B141, 253.

29. Arai I. and Fujii H. Proc. of the 19-th Intern. Conf. on High Energy Physics, Tokyo, 1978.

30. Feynman R.P., Kislinger M., Ravndal F. Phys.Rev. 1971, D3, 2706.

31. Kubota T., Ohta K. Phys.Lett. 1976, 65B, 374; Ohta K. Phys.Rev.Lett. 1979, 43, 1201.

32. Koniuk R., Isgur N. Phys.Rev. 1980, D21, 1868.

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РЕЛЯТИВИСТСКАЯ МОДЕЛЬ КВАРКОВ В СИСТЕМЕ БЕСКОНЕЧНОГО ИМПУЛЬСА И СТАТИЧЕСКИЕ СВОЙСТВА АДРОНОВ

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