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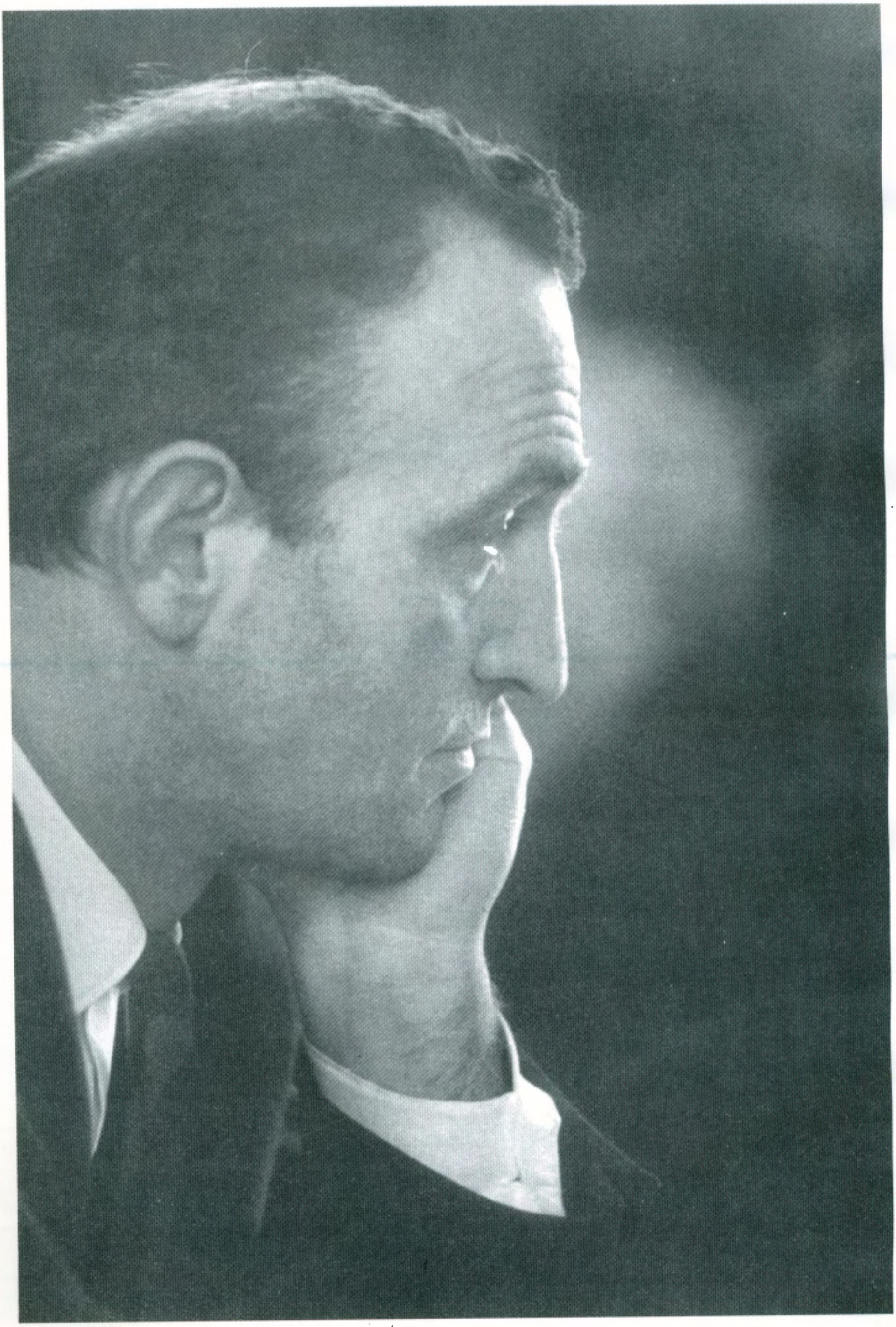

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ПРЕДИСЛОВИЕ

Настоящий сборник подготовлен к 70-летию со дня рождения академика Альберта Никифоровича Тавхелидзе, физика-теоретика с мировым именем. Сборник содержит краткий очерк жизни и деятельности ученого, полный список его научных трудов с классификацией по направлениям. Из этих многочисленных публикаций в сборнике воспроизводятся две лекции, прочитанные А.Н. Тавхелидзе 35 лет назад на знаменитом Семинаре по физике высоких энергий и элементарных частиц (Триест, Международный центр теоретической физики, 3 мая–30 июня 1965 г.), куда он был приглашен директором центра профессором Абдусом Саламом. В них впервые был изложен динамический подход Боголюбова–Тавхелидзе к построению адронов из *квазисвободных кварков*, обладающих новым квантовым числом — *цветом*, лежащим в основе современной квантовой хромодинамики.

СОДЕРЖАНИЕ

HIGH-ENERGY PHYSICS AND ELEMENTARY PARTICLES

LECTURES PRESENTED AT A SEMINAR, TRIESTE, 3 MAY–30 JUNE 1965, ORGANIZED BY THE

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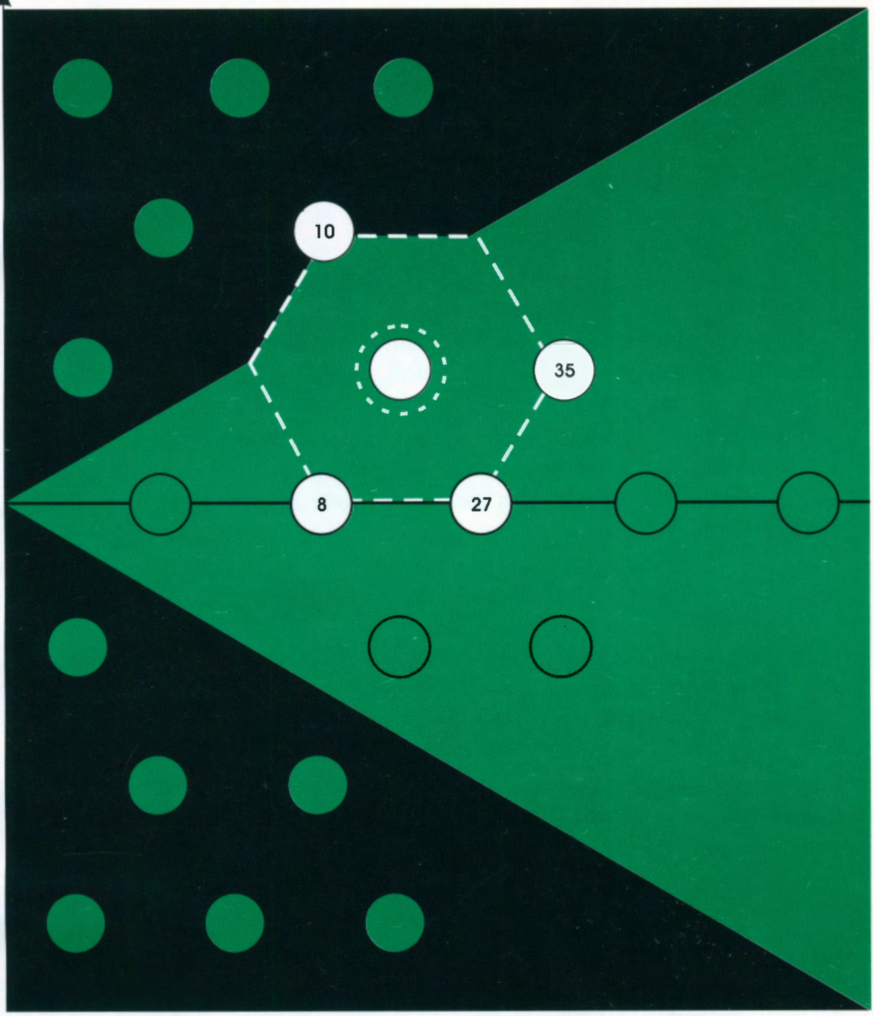

INTERNATIONAL ATOMIC ENERGY AGENCY, VIENNA, 1965

Lecture One

HIGHER SYMMETRIES AND COMPOSITE MODELS OF ELEMENTARY PARTICLES

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1. INTRODUCTION

Recently, group theoretical methods based on the synthesis of internal symmetry and space-symmetry properties have been extremely successful in the investigation of elementary particles properties.

There are two points of view in the application of group theoretical methods in the theory of elementary particles. According to the first one, all particles are considered as equally elementary. Another point of view suggests that there are some fundamental particles carrying all symmetry properties (as quarks, trions, Schwinger model, etc.), and all physical particles are considered as bound states of these. The composite model, in the $SU(6)$ frame, is much like the Wigner model of nuclear shells, combining spin independence and isotopic invariance of nuclear forces and classifying all nuclear states by different representations of the $SU(4)$ group. The difference consists in that we are not sure that particles representing the fundamental representation of $SU(3)$ exist at all. However, the plausibility of the idea of composite models of elementary particles attracts attention of many physicists.

Our aim is to discuss the simplest dynamical model and to investigate in the framework of this model the form factors and magnetic moments of baryons and mesons. This paper is in fact a review of investigations performed by Bogolubov et al. [1]. In order to illustrate the nature of the problems mentioned above, and as an introduction to the discussion of a relativistic model of composite particles, we shall first consider a nonrelativistic model, where the particles are regarded as bound states of the quarks.

Such a consideration may be justified by the following arguments. Examination of the relations predicted by $SU(6)$ suggests that the interaction, leading to formation of bound states, weakly depends on spin, and a good agreement of mass formulae with experimental data confirms an idea that masses of fundamental particles are large compared to the symmetry-breaking perturbation.

We shall show also that the small Dirac magnetic moment of the quark nevertheless could lead (due to the enhancement mechanism) to large moments of known baryons.

In the next part of our review we shall consider relativistic models of composite particles and their electromagnetic form factors.

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2. BOSONS

Consider a nonrelativistic model where the particles are regarded as bound states of the quarks. The wave functions of the quarks are the fundamental six-dimensional representation (sixtets) of $SU(6)$. It is well known that the members of the sixtet are $SU(3)$ triplets, each member having spin 1/2.

We denote the wave function of a quark by $\psi_a = \psi_{i,A}$ where $i = 1, 2$ is a spin state index and $A = 1, 2, 3$ is the $SU(3)$ index. For zero energy quark we have

$$\psi_a = t_A \chi_i$$
$$\chi_1 = \binom{1}{0} \quad \chi_2 = \binom{0}{1}, \tag{2.1}$$

t_A denotes the $SU(3)$-triplet. Thus for each A we have a two-component function ψ_a.

The transformation property of antiquark wave function is given by a conjugated representation of $SU(3)$ group. We denote the wave function of an antiquark by ψ^b, and for zero energy we have $\psi^{i,A} = \chi^i t^A \chi^i = \epsilon^{iK} \chi_K$.

Regarding a meson as a bound state of quark and antiquark, we denote the wave function of meson by ψ_a^b. The problem is to write down the equation for ψ_b^a.

To get such an equation we use the following assumption:

(a) we assume that the interaction which leads to a formation of bound states is weakly dependent on spin and unitary spin;

(b) the internal motion of a quark is regarded as nonrelativistic.

Taking into account the remarks which we made above, in zero approximation, where spin and unitary spin dependence is neglected, for the mesons we may write down the quasipotential equation which in the centre-of-mass system has the following form

$$(E^2 - H_0(r))\psi_a^b(r) = 0. \tag{2.2}$$

H_0 depends only on the relative co-ordinate $\vec{r}^{\,2} = (\vec{r}_1 - \vec{r}_2)^2$.

As H_0 is a scalar function, the energy spectrum for all 36 mesons is the same. In general, the perturbation which removes this degeneracy may depend on spin and unitary spin simultaneously. We assume that there are two types of perturbations: the first removes the degeneracy with respect to spin and the other with respect to unitary spin. We neglect the interference terms; they are supposed to be of second order in the perturbation.

The perturbation term which depends on spin may be represented in the following form

$$H_1 \psi_{iA}^{jB} = [(\vec{S}_i + \vec{S}_j)^2 V_1(r) + V_1'(r)]\psi_{iA}^{jB} + [(\vec{S}_i + \vec{S}_j)^2 V_2(r) + V_2'(r)]\psi_{iA}^{jA}. \tag{2.3}$$

The operators $\vec{S}_i$, $\vec{S}_j$ are the spin operators for the quark and antiquark, correspondingly. We use the designation

$$\begin{aligned} \vec{S}_i \psi_{iA}^{jB} &\equiv \vec{S}_i^{i'} \psi_{i'A}^{jB}, \\ \vec{S}_j \psi_{iA}^{jB} &= \vec{S}_{j'}^{j} \psi_{iA}^{j'B}. \end{aligned} \tag{2.4}$$

In order to write down a perturbation eliminating the unitary spin degeneracy we use the main idea of Zweig [2], viz., we assume that the mass of the quark t_3 is different from those of t_1

and t_2:

$$m_A^2 = m_0^2 + \Delta m^2 \delta_A^3 = \left(m_0^2 + \frac{1}{3}\Delta m^2\right) - (\lambda_8)_A \Delta m^2.$$

The perturbation which could give such a splitting may be chosen by analogy with Eq.(2.5)

$$H_2 \psi_a^b = W(r)[(\lambda_8)_A + (\lambda_8)_B]\psi_a^b. \tag{2.5}$$

The matrices $(\lambda_8)_A$ act on the index A and $(\lambda_8)_B$ on the index B:

$$\begin{aligned} (\lambda_8)_A \psi_a^b &= (\lambda_8)_A^{A'} \psi_{iA'}^{jB}, \\ (\lambda_8)_B \psi_a^b &= (\lambda_8)_{B'}^{B} \psi_{iA}^{jB'}. \end{aligned} \tag{2.6}$$

Taking into account the perturbation terms (2.3) and (2.4), we get

$$\begin{aligned} (E^2 - H_0)\psi_a^b &= [V_1(r)(\vec{S}_i + \vec{S}_j)^2 + V_1'(r)]\psi_{iA}^{jB} + \\ &+ [V_2(r)(\vec{S}_i + \vec{S}_j)^2 + V_2'(r)]\psi_{iA}^{jA} + \\ &+ W(r)[(\lambda_8)_A + (\lambda_8)_B]\psi_{iA}^{jB}. \end{aligned} \tag{2.7}$$

For the unperturbed solution we have

$$\psi_{iA}^{jB}(r) = \rho_{iA}^{jB}\psi(r), \tag{2.8}$$

where ρ_{iA}^{jB} is independent of r. After separating the spin and unitary spin structures we get

$$\rho_{iA}^{jB} = \begin{cases} V_A^B \chi_i^j & s = 1, \\ \rho_A^B \delta_i^j & s = 0, \end{cases} \tag{2.9}$$

s is the total spin of the system.

Since the perturbation terms conserve the total spin, Eq.(2.7) splits into two equations: one for spin 1 and the other for spin 0.

For the case of spin 1, we have 9 vector mesons:

$$V = \begin{pmatrix} \frac{1}{\sqrt{2}}(\rho_0 + \omega') & \rho^+ & K^{*+} \\ \rho^- & \frac{1}{\sqrt{2}}(\rho_0 - \omega') & \bar{K}^{*0} \\ K^{*-} & K^{*0} & \phi' \end{pmatrix}$$

$$V_1^2 = \rho^+, \quad V_2^1 = \rho^- \ldots \tag{2.10}$$

Using Eqs.(2.8) and (2.9) we see that to find the mass corrections to the vector mesons we must diagonalize the quadratic form

$$M^2 = \left(m_1^2 - \frac{a}{3}\right)\bar{V}_B^A \bar{V}_A^B + \frac{h_1}{3}\bar{V}_B^B \bar{V}_A^A + a(\bar{V}_3^A V_A^3 + \bar{V}_A^3 V_3^A). \tag{2.11}$$

Substituting Eq.(2.10) into Eq.(2.11) we get

$$\begin{aligned}
M^2 &= \left(m_1^2 - \frac{a}{3}\right)[\bar{\rho}^+\rho^+ + \bar{\rho}^-\rho^- \bar{\rho}^0\rho^0] + \\
&+ \left(m_1^2 + 2\frac{a}{3}\right)[\bar{K}^{*+}K^{*+} + \bar{K}^{*-}K^{*-} + \bar{K}^{*0}K^{*0} + K^{*0}\bar{K}^{*0}] + \\
&+ \left(m_1^2 - \frac{a}{3}\right)\bar{\omega}'\omega' + \left(m_1^2 - \frac{5}{3}a\right)\bar{\phi}'\phi' + \\
&+ \frac{h_1}{3}(2\bar{\omega}'\omega' + \bar{\phi}'\phi' + \sqrt{2}\bar{\omega}'\omega' + \sqrt{2}\bar{\phi}'\phi'). \qquad (2.12)
\end{aligned}$$

From this we see that

$$m_\rho^2 = m_1^2 - \frac{a}{3}, \quad m_{K^*}^2 = \left(m_1^2 + \frac{2}{3}a\right). \qquad (2.13)$$

If we put $h_1 = 0$, we get

$$m_{\omega'}^2 = m_\rho^2 \quad \text{and} \quad m_{\phi'}^2 + m_\rho^2 = 2m_{K^*}^2. \qquad (2.14)$$

If we want to calculate the masses of real ω and ϕ mesons, we must diagonalize the quadratic form

$$\left(m_1^2 - \frac{a}{3}\right)\bar{\omega}'\omega' \left(m_1^2 - \frac{5}{3}a\right)\bar{\phi}'\phi' + \frac{h_1}{3}[2\bar{\omega}'\omega' + \bar{\phi}'\phi' + \sqrt{2}\bar{\omega}'\phi' + \sqrt{2}\bar{\phi}'\omega']. \qquad (2.15)$$

The eigenvalues of this form are given by the equation

$$\begin{vmatrix} m_1^2 - \frac{a}{3} + \frac{2h_1}{3} - \mu & \frac{\sqrt{2}}{3}h_1 \\ \frac{\sqrt{2}}{3}h_1 & m_1^2 - \frac{5}{3}a + \frac{h_1}{3} - \mu \end{vmatrix} = 0. \qquad (2.16)$$

Excluding the parameters m_1^2, a, and h_1 we obtain the well-known formula of Schwinger [3].

$$(m_\omega^2 - m_\rho^2)(m_\phi^2 - m_\rho^2) = \frac{4}{3}(m_{K^*}^2 - m_\rho^2)(m_\phi^2 + m_\omega^2 - 2m_{K^*}^2). \qquad (2.17)$$

Now we would like to give a brief survey of the situation concerning pseudoscalar mesons. To get the known mass formula it is sufficient to make replacements $m_1 \to m_0$, $h_1 \to h_0$, $\rho \to \pi$, $K \to K^*$, $\omega' \to \eta'$, $\phi' \to \eta''$. Then we have

$$m_\pi^2 = m_0^2 - \frac{a}{3}; \quad m_K^2 = m_0^2 + \frac{2}{3}a. \qquad (2.18)$$

From this follows the Glashow formula

$$m_{K^*}^2 = m_\rho^2 = m_K^2 - m_\pi^2. \qquad (2.19)$$

In order to avoid the mixing of unitary singlet and octet of pseudoscalar mesons, we choose from 36-meson states only 35 states which satisfy the conditions

$$\psi_a^a = 0. \qquad (2.20)$$

We also require that the perturbation does not break this condition. For the perturbation matrix Γ this requirement leads to

$$\Gamma^{c,a}_{c,b} = 0. \tag{2.21}$$

This condition may be satisfied if we take the perturbation matrix in the following form

$$\Gamma^{iA;i'A'}_{jB;j'B'} = g^{iA;i'A'}_{jB;j'B'} - \frac{1}{6}\delta^A_B\delta^j_i g^{kc;i'A'}_{kc;j'B'}.$$

Then for the pseudoscalar octet we get

$$m^2_\eta = m^2_0 + a$$

which leads us to the Gell-Mann–Okubo mass formula

$$m^2_\eta = \frac{4}{3}m^2_K - \frac{1}{2}m^2_\pi.$$

3. BARYONS

In our approach baryons are regarded as bound states of three quarks. Denote the wave function of the baryon by ψ_{abc}. In the zero approximation we write the following equation for baryons

$$(E - H_0)\psi_{abc} = 0. \tag{3.1}$$

We represent the perturbation terms in the following form

$$\begin{aligned} H\psi_{abc} &= [(\vec{S}_i + \vec{S}_j + \vec{S}_K)^2 W_1 + W_2']\psi_{iA,jB,kC} \\ &+ W_2[(\lambda_8)_A + (\lambda_8)_B + (\lambda_8)_C]\psi_{abc} \\ &+ W_3[(\lambda_8)_A P_{BC} + (\lambda_8)_B P_{AC} + (\lambda_8)_C P_{AB}]\psi_{iA,jB,kC}, \end{aligned} \tag{3.2}$$

the functions W depend only on a relative distance between the quarks, P_{AB} is the permutation operator. The third term carries on the exchange interaction for the octet.

There are three types of symmetries for the wave function with respect to the permutation of the unitary and spin indices: completely symmetric with 56 components, antisymmetric with 20 components and the function transforming under the two-dimensional representation of the permutation groups of three elements, 70 components.

As was shown in paper [4], a better relation between the magnetic moments is obtained for the function symmetric with respect to spin and unitary spin, 56 components.

However, if we assume that there are no additional quantum numbers, the space function must be completely antisymmetric which is impossible if the system is in the S-state.

A way out can be found if the quarks are supposed to obey parastatistics [5] or if we introduce additional quantum numbers which antisymmetrize the total wave function. Employing these additional quantum numbers we are able to make the quark charges integer without violating the relations between the magnetic moments.

Consider Eq.(3.1) with the perturbation (3.2) for the functions symmetric with respect to spin and unitary spin.

The perturbation term (3.2) conserves the total spin of the system, and the equation for the baryons splits into two equations: one — for spin 3/2 and the other — for spin 1/2. In the zero approximation, the baryons wave function can be represented in the form

$$\psi_{ABC}(r_1, r_2, r_3) = \varphi(r_1, r_2, r_3)\phi_{ABC}. \tag{3.3}$$

For the octet particles we have $\phi_{ai,bj,ck}$

$$\begin{aligned}
p &= \frac{1}{\sqrt{3}}\{\phi_{11,11,22} - \phi_{11,21,12}\}, \\
r &= \frac{1}{\sqrt{3}}\{\phi_{21,21,22} - \phi_{21,11,22}\}, \\
\Sigma^+ &= \frac{1}{\sqrt{3}}\{\phi_{11,11,32} - \phi_{11,31,12}\}, \\
\Sigma^- &= \frac{1}{\sqrt{3}}\{\phi_{21,21,32} - \phi_{21,31,22}\}, \\
\Sigma^0 &= \frac{1}{\sqrt{6}}\{2\phi_{11,21,32} - \phi_{11,31,22} - \phi_{21,31,12}\}, \\
\Lambda &= \frac{1}{\sqrt{2}}\{\phi_{11,31,22} - \phi_{21,31,12}\}, \\
\Xi^- &= \frac{1}{\sqrt{3}}\{\phi_{31,31,22} - \phi_{21,31,32}\}, \ \Xi^0 = \frac{1}{\sqrt{3}}\{\phi_{31,31,12} - \phi_{11,31,32}\}.
\end{aligned} \tag{3.4}$$

The functions ϕ_{ABC} are normalized by the condition

$$||\phi_{AAA}||^2 = 6, \quad ||\phi_{AAB}||^2 = 2, \quad ||\phi_{ABC}|| = 1.$$

Using the perturbation theory we obtain the following relations for the masses of baryons

$$m_\Sigma - m_N = a + b, \quad m_{\Xi^-} - m_N = 2a - b, \quad m_\Lambda - m_N = a - b. \tag{3.5}$$

From this, follows the Gell-Mann–Okubo mass formula

$$m_\Xi + m_N = \frac{3m_\Lambda + m_\Sigma}{2}. \tag{3.6}$$

For the decuplet particles we have

$$m_0 - m_{\Sigma^*} = m_{\Sigma^*} - m_{\Xi^*} = m_{\Xi^*} - m_{\Omega^-} = -a - b. \tag{3.7}$$

Thus, for the mass splitting in the octet and decuplet we get

$$m_\Sigma - m_N = m_{\Sigma^*} - m_0. \tag{3.8}$$

It is easy to write the perturbation which describes the electromagnetic mass splitting within the isotopic multiplets

$$\begin{aligned}
H_2\psi_{abc} &= \{I_1(e_A^2 + e_B^2 + e_C^2) + I_2(e_A e_B + e_A e_C + e_B e_C) \\
&+ I_3(e_A e_B P_{AB} + e_A e_C P_{AC} + e_B e_C P_{BC}) e\psi_{abc}\}.
\end{aligned} \tag{3.9}$$

The first term represents the electromagnetic corrections to the proper effective mass of the quarks; the second term, the direct electromagnetic interaction of the quarks; and the third term, the exchange interaction. Noting that $e_A^2 = \frac{2}{3} + \frac{1}{3}e_A$ we write down the electromagnetic perturbation as

$$\begin{aligned} H_2 &= \{I_1'(e_A + e_B + e_C) + I_2'(e_A + e_B + e_C)^2 \\ &+ I_3(e_A e_B P_{AB} + e_A e_C P_{AC} + e_B e_C P_{BC})\} \psi_{abc}. \end{aligned} \tag{3.10}$$

Using this interaction we get

$$\begin{aligned} m_p - m_n &= \alpha + \beta + \frac{1}{3}\gamma, \quad m_{\Sigma^+} - m_{\Sigma^-} = 2\alpha + \frac{2}{3}\gamma, \\ m_{\Xi^-} - m_{\Xi^0} &= -\alpha + \beta - \frac{1}{3}\gamma, \\ m_{\Delta^+} - m_{\Delta^0} &= m_{\Sigma^{*+}} - m_{\Sigma^{*0}} = \alpha + \beta + \frac{1}{3}\gamma, \\ m_{\Delta^-} - m_{\Delta^0} &= m_{\Xi^{*-}} - m_{\Xi^{*0}} = m_{\Sigma^{*-}} - m_{\Sigma^{*0}} = -\alpha + \beta + \frac{2}{3}\gamma, \\ m_{\Delta^{++}} - m_{\Delta^+} &= \alpha + 3\beta + \frac{4}{3}\gamma. \end{aligned} \tag{3.11}$$

From this we find the electromagnetic mass splitting in the octet

$$m_p - m_n + m_{\Sigma^-} - m_{\Sigma^+} = m_{\Xi^-} - m_{\Xi^0}. \tag{3.12}$$

We also get the relations between the electromagnetic mass splitting in octet and decuplet

$$m_p - m_n = -m_{\Delta^+} - m_{\Delta^0} = -m_{\Delta^-} - m_{\Delta^0} = m_{\Xi^{*+}} - m_{\Xi^{*0}}. \tag{3.13}$$

From $SU(6)$ symmetry follow relations between the magnetic moments of baryons. This is one of the important consequences of $SU(6)$ symmetry. Using the formulae of quantum mechanics for the relativistic composite model it is possible to calculate the magnetic moments of baryons.

Actually, choosing the magnetic moment operator in the form

$$\vec{\mu} = 3\mu[(e\vec{\sigma})_A + (e\vec{\sigma})_B + (e\vec{\sigma})_C] \tag{3.14}$$

one can easily show that the magnetic moments of the resonances are $3e\mu$ where e is the resonance charge, and for the baryon octet we have

$$\mu_p = 3\mu, \quad \mu_m = 2\mu, \quad \mu_{\Sigma^-} = -\mu. \tag{3.15}$$

The magnetic moments of other particles follow from the unitary symmetry relations

$$\mu_p = \mu_{\Sigma^+}, \quad \mu_n = 2\mu_{\Xi^0} = 2\mu_\Lambda = -2\mu_{\Sigma^0}. \tag{3.16}$$

In order to get the experimental value of the magnetic moment we must put μ equal to the nuclear magneton $\mu = e/2m$, where m is the mass of the bound states, that is, baryons.

Now we want to study a mechanism providing for such an enhancement of the magnetic moment of the quark.

Consider the bound states of the Dirac particle in the scalar field. We wish to calculate the magnetic moment of Dirac particles which are in the bound states.

Let us write the Dirac equation

$$[M + V(r) + \beta\vec{\alpha} \cdot (\vec{p} - e\vec{A})]u = E\beta u. \tag{3.17}$$

We consider the case

$$V(r) = \begin{cases} 0 & r > r_0 \\ -M + \alpha & r < r_0 \end{cases}. \tag{3.18}$$

Representing the spinor u in the form

$$u = \begin{pmatrix} \varphi \\ \chi \end{pmatrix} \tag{3.19}$$

we get the following equation for φ:

$$(M + V - E)\varphi + \vec{\sigma}(\vec{p} - e\vec{A})\frac{1}{M + E + V}\vec{\sigma}(\vec{p} - e\vec{A})\varphi = 0. \tag{3.20}$$

Taking into account only the first order terms in the electromagnetic interaction, we obtain the correction to the energy operator

$$\Delta\hat{E} = e\vec{\sigma} \cdot \vec{A}\frac{1}{M + E + V}\vec{\sigma} \cdot \vec{p} + e\vec{\sigma} \cdot \vec{p}\frac{1}{M + E + V}\vec{\sigma} \cdot \vec{A}. \tag{3.21}$$

Calculating the energy spectrum correction to the S-state, we arrive at the following expression for the magnetic moment

$$\mu = \frac{2e}{3}\int_0^\infty \frac{g(r)}{M + E + V}\frac{dg}{dr}r^3\,dr, \tag{3.22}$$

where the function $W = rg$ satisfies the equation

$$W'' + (E^2 - M - V)W = 0. \tag{3.23}$$

Using the requirement that $W(r_0 - 0) = W(r_0 + 0)$ we obtain the equation for the energy spectrum of bound states

$$\frac{E - M}{r_0(E + M)(E + \alpha)} + \frac{1}{E + \alpha}\left(\frac{W'}{W}\right)_{r_0 - 0} = \frac{1}{E + M}\left(\frac{W'}{W}\right)_{r_0 + 0}. \tag{3.24}$$

Later on, the calculations are made for the case $M \to \infty$. Here the energy spectrum equation has the form

$$r_0(E^2 - \alpha^2)^{1/2}\operatorname{cotg} r_0(E^2 - \alpha^2)^{1/2} = 1 - r_0(E + \alpha). \tag{3.25}$$

Integrating Eq.(3.22) and taking the limit $M \to \infty$, we get

$$\mu = \frac{e}{6}\frac{4Er_0 + 2\alpha r_0 - 3}{2E^2r_0 - 2E_0 + \alpha} \tag{3.26}$$

for the case $\alpha = 0$, numerical calculations give $Er_0 \sim 2$. Thus,

$$\mu \sim \frac{e}{2E}\frac{5}{6}. \tag{3.27}$$

If the proton is considered as a bound state of three quarks in the self-consistent field, the energy of each quark is equal to $E = m/3$. Using then (3. 27) for the magnetic moment of the proton we obtain a value for the nuclear magneton to be $\mu_Q \sim 2.5$. We have seen that the effective magnetic moment of Dirac particles which are in the bound states is determined by the bound state energy.

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Lecture Two

ELECTROMAGNETIC FORM FACTORS IN COMPOSITE MODELS OF ELEMENTARY PARTICLES (RELATIVISTIC MODELS)

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1. INTRODUCTION

In our previous review we have considered the simplest variant of the dynamical approach to $SU(6)$ symmetry. Using the assumption that the mass of a quark is large we write down nonrelativistic $SU(6)$ symmetric equations which reflect mass relations obtained in $SU(6)$ and relations between magnetic moments. In the simplest example of a Dirac particle scalar field, the mechanism of the enhancement of the magnetic moment was shown.

It is of interest to write down relativistic invariant equations for composite particles which describe the enhancement of magnetic moments and produce the true electromagnetic form factors of elementary particles**.

Recently a number of very interesting results concerning form factors were given in the theory of Salam et al. [1] based on $S\tilde{U}(12)$ symmetry. However, in order to calculate electromagnetic form factors, a nonminimal electromagnetic interaction is introduced there.

In our papers**, we studied a dynamical model of composite particles, satisfying requirements of $\tilde{U}(12)$ invariance as well as relativistic ones. In this model the minimal electromagnetic interaction for the quarks is suggested, and electromagnetic form factors and decay vertices for baryons and bosons are obtained.

2. RELATIVISTIC EQUATION FOR MESONS [1]

We consider the composite model of elementary particles in which all the particles are regarded as bound states of three basic particles with spin 1/2. In a relativistic extension of the $SU(6)$ symmetry theory, the quarks must be described by 12-component functions $\psi_A \equiv \psi_{\alpha,a}$, where $a = 1,2,3$ is the $SU(3)$ index, and $\alpha = 1,2,3,4$ is the Dirac spinor index. The group $SU(12)$ acting on the quark function ψ is generated by 144 matrices $\lambda_i\Gamma_r$, where λ_i are the nine generators of U(3), and Γ_r are the sixteen Dirac matrices which form $\tilde{U}(4)$. We denote the antiquark wave function by $\psi^B \equiv \psi^{\beta,b}$.

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**See Ref. [1] of Tavkhelidze, A., Higher Symmetries and Composite Models of Elementary Particles, these Proceedings.

Mesons which are bound states of the quark and antiquark are described by a mixed spinor of the second rank ψ_A^B. The question is to find an equation for $\psi_A^B(x_1, x_2)$.

In order to study the bound states of the system one could resort to the methods employed in quantum field theory, viz., to the Bethe–Salpeter equation. However, there are well-known difficulties connected with a solution of this equation in the case of strong interactions.

We will therefore start from the relativistic-invariant equation for two particles admitting solutions of the form

$$\psi_A^B(x_1, x_2) = \phi_1(x_1 - x_2)\phi_A^B(x_1 + x_2), \tag{2.1}$$

where $\phi_A^B(x_1 + x_2)$ is a mixed spinor of the second rank which describes the motion of the system as a whole, and $\phi(x_1 - x_2)$ is a scalar function describing the relative motion of particles. The simplest example of an equation invariant under a homogeneous group* $\tilde{U}(12)$ and describing the interaction of two particles is the squared Dirac equation with the factorizing potential:

$$\begin{aligned} \mathcal{D}_A^{A'}(x_1)\bar{\mathcal{D}}_{B'}^{B}(x_2)\psi_{A'}^{B'}(x_1, x_2) &= -igW(x_1 - x_2)\times \\ \times \int dx_1' dx_2' W(x_1' - x_2')\delta(x_1 + x_2 - x_1' - x_2')\psi_A^B(x_1', x_2'), \end{aligned} \tag{2.2}$$

where $\mathcal{D}(\bar{\mathcal{D}})$ is the squared Dirac operator for the particle (antiparticle)

$$\mathcal{D}_{A(B')}^{A'(B)}(x) = \left[\left(M \mp i\sum_{n=0}^{3}\gamma^n\frac{\partial}{\partial x^n}\right)\right]\left[\left(M \pm i\sum_{n=0}^{3}\gamma^n\frac{\partial}{\partial x^n}\right)\right] \tag{2.3}$$

and

$$\gamma^n\gamma^m + \gamma^m\gamma^n = 2g^{mn}, \quad g^{11} = g^{22} = g^{33} = -g^{00} = -1.$$

We require that the Fourier transform $W(q^2)$ of the scalar function $W(x)$ be real in the Euclidean space. In particular, $W(q^2)$ can be taken in the form

$$W(q^2) = \sum_N \frac{C_N}{(q^2 - N^2 - i\epsilon)^2}. \tag{2.4}$$

The operators $\mathcal{D}$ and W do not act on the spinor and unitary indices. Thus, Eq.(2.2) is invariant under $\tilde{U}(12)$ in a trivial fashion. Before including into Eq.(2.2) the symmetry breaking terms we shall study it in more detail.

Performing the Fourier transform

$$\begin{aligned} \psi_A^B(x_1, x_2) &= \int \psi_A^B(p_1, p_2)e^{-ip_1x_1 - ip_2x_2}dp_1\, dp_2 \\ W(x_1 - x_2) &= \int W(q^2)e^{-iq(x_1 - x_2)}dq \end{aligned} \tag{2.5}$$

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*$\tilde{U}(12)$ acts only on the discrete indices and does not affect the variables x.

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we write Eq. (2.2) in the form

$$(p_1^2 - M^2)(p_2^2 - M^2)\psi_A^B(p_1, p_2) = -ig_0\pi^4 W(q^2)\phi_A^B(p)$$
$$p = p_1 + p_2 \quad q = \tfrac{1}{2}(p_1 - p_2), \tag{2.6}$$

where

$$\phi_A^B(p) = \int W(q^2)\psi_A^B\left(\frac{p}{2} + q, \ \frac{p}{2} - q\right). \tag{2.7}$$

Assuming the quark mass M to be large (infinite in the limit) and taking into account only the higher terms in M, we obtain in the zero approximation

$$\psi_A^B\left(\frac{p}{2} + q, \ \frac{p}{2} - q\right) = -igW(q^2)\phi_A^B(p)$$
$$g = \frac{g_0\pi^4}{M^4}. \tag{2.8}$$

Multiplying (2.8) by $W(q^2)$ and integrating over q we get

$$1 = -ig\int W^2(q^2)dq. \tag{2.9}$$

The function $\phi_A^B(p)$ describing the motion of the system as a whole remains undetermined. Going over to the Euclidean momenta $q_0 \to iq_4$ and taking into account (2.9) we obtain the real value for the charge

$$1 = g\int W^2(-q_E^2)dq_E, \tag{2.10}$$

q_E is the four-momentum in the Euclidean space.

In order to get equations for the function $\phi_A^B(p)$ we rewrite (2.6) as

$$\left(\frac{1}{2}p^2 + 2q^2\right)\psi_A^B\left(\frac{p}{2} + q, \ \frac{p}{2} - q\right) = \frac{1}{M^2}(M^4 + p_1^2p_2^2)\psi_A^B\left(\frac{p}{2} + q, \ \frac{p}{2} - q\right)$$
$$+ \ i\frac{g_0\pi^4}{M^2}W(q^2)\phi_A^B(p). \tag{2.11}$$

Multiplying Eq. (2.11) by $W(q^2)$, integrating over q and using Eqs.(2.8) and (2.9) we get

$$(p^2 - m^2)\phi_A^B(p) = 0, \tag{2.12}$$

where the meson mass is given by the expression

$$m^2 = 4ig\int q^2W^2(q^2)dq = 4\frac{\int q_E^2 W^2(q_E^2)dq_E}{\int W^2(-q_E^2)dq_E}. \tag{2.13}$$

Later on we use the solution of Eq.(2.12) which satisfies the supplementary condition:

$$(\gamma_0 \cdot p)_A^{A'}\phi_{A'}^B(p) + (\gamma \cdot p)_{B'}^B\phi_A^{B'}(p) = 0. \tag{2.14}$$

This condition chooses the solutions corresponding to opposite signs of the energies for the quark and antiquark. This can be easily seen if we rewrite (2.14) in the rest frame:

$$(\gamma_0)_A^{A'}\phi_{A'}^B + (\gamma_0)_{B'}^B\phi_A^{B'} = 0. \tag{2.15}$$

Since $p^2 = m^2$, we can require a simultaneous fulfilment of the Eqs. (a) or (b):

$$\begin{aligned} &(\hat{p}-m)_A^{A'}\phi_{A'}^B = 0, \quad (\hat{p}+m)_{B'}^B\phi_A^{B'} = 0, \quad (a) \\ &(\hat{p}+m)_A^{A'}\phi_{A'}^B = 0, \quad (\hat{p}-m)_{B'}^B\phi_A^{B'} = 0. \quad (b) \end{aligned} \tag{2.16}$$

Further we shall use Eqs.(2.16a). Note that the supplementary condition (2.16) breaks up the $\tilde{U}(12)$ invariance. Indeed, in the rest frame $\vec{p} = 0$ (2.16a) reduces to (2.14) which is invariant under the group $SU(6)$.

We note that the condition (2.14) selects the mesons with a negative intrinsic parity.

Let us investigate in more detail the spinor and unitary structure of the functions ϕ_A^B. The mixed spinor of the second rank can be represented as

$$\phi_{\alpha,p}^{\beta,q} = [\phi^i + \gamma^5\phi_5^i + i\gamma^\mu\gamma^5\phi_{\mu 5}^i + \gamma^\mu\phi_\mu^i + \frac{1}{2}\sigma^{\mu\nu}\phi_{\mu\nu}^i]_\alpha^\beta(\lambda_i)_p^q \tag{2.17}$$

where $i = 0\ldots 8$, $\sigma^{\mu\nu} = \frac{i}{2}(\gamma^\mu\gamma^\nu - \gamma^\nu\gamma^\mu)$ and λ_i are Gell-Mann matrices.

Substituting Eq. (2.17) into Eqs. (2.16) we get [2]

$$\begin{aligned} &p_\mu\phi_5^i = im\phi_{\mu 5}^i, \quad \phi^i = 0, \quad p_\mu\phi_\nu^i - p_\nu\phi_\mu^i = im\phi_{\mu\nu}^i, \\ &p_\mu\phi_{\mu 5}^i = -im\phi_5^i, \qquad\qquad p^\nu\phi_{\nu\mu}^i = -im\phi_\mu^i. \end{aligned} \tag{2.18}$$

It follows from here the vector field should be transverse:

$$p_\mu\phi_\mu^i = 0$$

$$\phi_A^B = \left[\left(1 + \frac{\hat{p}}{m}\right)(\gamma_5\phi_5 + \gamma_\mu\phi_\mu)\right]_A^B.$$

3. FORM FACTORS AND THE MAGNETIC MOMENTS OF MESONS

To calculate the form factors of mesons, we consider Eqs.(2.2) in a weak external electromagnetic field. For this purpose we make the usual substitution

$$i\frac{\partial}{\partial x^m} \rightarrow i\frac{\partial}{\partial x^m} + eA_m(x) \tag{3.1}$$

and write Eq. (2.2) in the form

$$\begin{aligned} &\mathcal{D}_A^{A'}(x_1, A)\bar{\mathcal{D}}_{B'}^B(x_2, A)\psi_{A'}^{B'}(x_1, x_2) = -ig_0W(x_1 - x_2)\times \\ &\times\int dx_1'dx_2'W(x_1' - x_2')\delta(x_1 + x_2 - x_1' - x_2')\psi_A^B(x_1', x_2'), \end{aligned} \tag{3.2}$$

where

$$\mathcal{D}_A^{A'}(x,A) = \left[\left(M - i\gamma^n \frac{\partial}{\partial x^n} - e\hat{A}(x)\right)\left(M + i\gamma^n \frac{\partial}{\partial x^n} + e\hat{A}(x)\right)\right]_A^{A'} \tag{3.3}$$

$$\mathcal{D}_{B'}^{B}(x,A) = \left[\left(M - i\gamma^n \frac{\partial}{\partial x^n} + e\hat{A}(x)\right)\left(M + i\gamma^n \frac{\partial}{\partial x^n} + e\hat{A}(x)\right)\right]_{B'}^{B}$$

Here e is the matrix of the electric charge which is expressed in a certain way in terms of the matrices λ.

Retaining in Eq.(3.3) the first order terms in e and making use of the expansion in the large mass of the quark M, we get an equation describing the motion of the meson in a weak electromagnetic field

$$(p^2 - m^2)\phi_A^B(p) = 2\int dk\, f(k^2) p \cdot A(k)[e_A^{A'}\phi_{A'}^B + e_{B'}^B \phi_B^{B'}(p-k)] +$$

$$+2\int dk\, f(k^2)[(e\hat{k}\hat{A})_A^{A'}\phi_{A'}^B(p-k) + (e\hat{A}\hat{k})_{B'}^B \phi_A^{B'}(p-k)], \tag{3.4}$$

or, in the x-representation,

$$(\Box_x + m^2)\phi_A^B(x) = 2i\tilde{A}_\mu(x)\left[e_A^{A'}\frac{\partial}{\partial x_\mu}\phi_{A'}^B + e_{B'}^B \frac{\partial}{\partial x_\mu}\phi^{B'}(x)\right] +$$

$$+2i\left[\left(e\gamma^\mu\gamma^\nu \frac{\partial \tilde{A}_\nu}{\partial x^\mu}\right)_A^{A'}\phi_{A'}^B(x) - \left(e\gamma^\mu\gamma^\nu \frac{\partial \tilde{A}_\nu}{\partial x^\mu}\right)_{B'}^{B}\phi_A^{B'}(x)\right], \tag{3.4a}$$

$$\tilde{A}(x) = \int \mathrm{e}^{-ikx} f(k)A(k)dk,$$

where $f(k^2)$ is determined in the following way

$$f(k^2) = \frac{\int dq\, W\left(q - \frac{k}{2}\right) W(q)}{\int dq\, W^2(q)}. \tag{3.5}$$

Before calculating the form factor of mesons we consider the case of the constant magnetic field. Using Eq.(2.16) one can see that for large components ϕ we obtain

$$\phi = a + \vec{\sigma} \cdot \vec{\varphi} \tag{3.6}$$

and for the vector field $\vec{\varphi}$ in the constant magnetic field the following equation holds

$$(p^2 - m^2)\vec{\varphi} = 2ie[\vec{H} \cdot \vec{\varphi}], \tag{3.7}$$

where e is the charge of the vector meson.

It follows from Eq. (3.7) that the magnetic moment of the vector meson is equal to e/m. We make two important remarks: (a) the expression for the magnetic moment involves not the quark mass M, but the meson mass m; (b) the magnetic moment of the meson regarded as a composite particle is twice as big as the normal magnetic moment of the vector meson with charge e and mass m. Now we pass to a calculation of the meson form factors.

Let us rewrite Eq.(3.4) in a more compact form

$$(p^2 - m^2)\phi_A^B = \int dk\, f(k^2) \left[\Gamma_A^{A'}(k,p)\phi_{A'}^B(p-k) + \Gamma_{B'}^B(p,k)\phi_A^{B'}(p-k)\right], \tag{3.8}$$

where the notations are introduced

$$\Gamma_A^{A'} = 2[p \cdot A(k)e + e\hat{k} \cdot \hat{A}]_A^{A'},$$
$$\Gamma_{B'}^B = 2[p \cdot A(k)e + e\hat{A} \cdot \hat{k}]_{B'}^B. \tag{3.9}$$

By definition the meson vertex is given by

$$\mathcal{J}(p, p-k) = 2mJ_\alpha \cdot A^\alpha \equiv \bar{\phi}_B^A(p)\Gamma_A^{A'}(p,k)\phi_{A'}^B(p-k) + \bar{\phi}_B^A\Gamma_{B'}^B(p,k)\phi_A^{B'}(p-k). \tag{3.10}$$

Substituting the spinor ϕ_B^A in the form of Eq. (2.17) and using Eq. (2.18) we get the electromagnetic vertex in the following form

$$\begin{aligned}J_\alpha = \frac{p_\alpha}{m} f(k^2) \Big\{ \left(1 + \frac{k^2}{8m^2}\right)(\bar{\varphi}_\mu\varphi_\mu)_F^e - \left(1 + \frac{k^2}{4m^2}\right)(\bar{\varphi}_5\varphi_5)_F^e - \\ -\frac{1}{2m^2}(k^\mu k^\nu - \frac{1}{4}g^{\mu\nu}k^2)(\bar{\varphi}_\mu\varphi_\nu)_F^e \Big\} + \frac{k_\mu}{m} f(k^2)[\bar{\varphi}_\mu\varphi_\alpha - \bar{\varphi}_\alpha\varphi_\mu]_F^e + \\ + \frac{1}{m^2}\epsilon^{\alpha\beta\sigma\rho}p_\sigma k_\rho[\bar{\varphi}_\beta\varphi_5 - \bar{\varphi}_5\varphi_\beta]_D^e,\end{aligned} \tag{3.11}$$

where

$$(\bar{\varphi}\varphi)_F^e = \bar{\varphi}_{q'}^p\varphi_q^{q'}e_p^q - \bar{\varphi}_{q'}^p e_q^{q'}\varphi_p^q,$$
$$(\bar{\varphi}\varphi)_A^e = \bar{\varphi}_{q'}^p\varphi_q^{q'}e_p^q + \bar{\varphi}_{q'}^p e_q^{q'}\varphi_p^q. \tag{3.12}$$

From this expression we have

$$g_E^V = ef(k^2)\left(1 + \frac{k^2}{8m^2}\right), \quad g_E^P = ef(k^2)\left(1 + \frac{k^2}{4m^2}\right), \tag{3.13}$$

$$g_Q^V = \frac{2e}{m^2} f(k^2), \quad g_M^V = \frac{e}{m} f(k^2),$$

where $g_E^{V,P}$ is the electric form factor for the vector and pseudoscalar mesons respectively, g_M^V is the magnetic form factor for vector mesons, and g_Q^V is the quadrupole moment of vector mesons.

The last term in (3.11) describes the radiative decay of vector mesons. Substituting there the known matrices for pseudoscalar and vector mesons, we get

$$[\bar{\varphi}_5\varphi_\beta]^e_D = \frac{1}{3}\left[\bar{\pi}^+\rho^+ + \bar{\pi}^-\rho^- + \bar{K}^+K^+ + \bar{K}^-K^- - 2\bar{K}^0K^{*0} - 2K^0\bar{K}^{*0}\right] +$$
$$+\left[\frac{1}{3}\bar{\pi}^0 + \frac{1}{\sqrt{3}}\bar{\eta}^0 + \sqrt{\frac{2}{3}}\bar{\chi}^0\right]\rho_0 + \left[\frac{1}{\sqrt{3}}\bar{\pi}^0 + \frac{1}{3}\bar{\eta}^0\right]\varphi^0 + \tag{3.14}$$
$$+\left[\sqrt{\frac{2}{3}}\bar{\pi}^0 + \frac{\sqrt{2}}{3}\bar{\eta}^0 + \frac{2}{3}\bar{\chi}^0\right]\omega^0.$$

For the physical particles ω and φ we have

$$\omega = \sqrt{\frac{2}{3}}\omega_0 + \sqrt{\frac{1}{3}}\varphi_0,$$
$$\varphi = \sqrt{\frac{1}{3}}\omega_0 - \sqrt{\frac{2}{3}}\varphi_0. \tag{3.15}$$

The matrix elements for the radiative decay are given by the expression

$$M = -\frac{2e}{m}g_{VP}\epsilon^{\alpha\beta\sigma\rho}p_\sigma k_\rho A_\alpha(\bar{\varphi}_5\varphi_\beta)^e_D, \tag{3.16}$$

where the radiative decay constants are equal to

$$g_{\omega\pi^0} = 3g_{\rho^0\pi^0} = \sqrt{3}g_{\rho^0\eta^0} = 3g_{K^{*+}K^+} = 3g_{K^{*-}K^-} = 3\sqrt{3}g_{\omega\eta^0} = \frac{3\sqrt{3}}{2\sqrt{2}}g_{\varphi\eta_0} = -\frac{3}{2}g_{K^{*0}K^0}. \tag{3.17}$$

Taking into account the mass difference between the mesons only in the phase space, we get the following expression for the decay width

$$\Gamma_{V\to p+\gamma} = \alpha\frac{4}{3}m_V\left(\frac{m_V^2 - m_p^2}{2m_V^2}\right)^3 g^2_{Vp}, \tag{3.18}$$

where α is the fine structure constant, m_V and m_p are the masses of vector and pseudoscalar mesons.

Now we write down the matrix element for the decay of a heavy pseudoscalar meson χ_0 into the vector mesons:

$$M = \frac{2e}{M}g_{\chi_0V}\epsilon^{\alpha\beta\sigma\rho}p_\sigma k_\rho A_\alpha\bar{\varphi}_\beta\varphi_5. \tag{3.19}$$

Here g_{χ_0V} is the radiative decay constant of χ_0 meson. We carry out calculations without mixing χ_0 and η_0 mesons. We obtain

$$g_{\chi_0\rho_0} = \frac{3}{2}g_{\chi_0\omega} = \sqrt{\frac{2}{3}} \tag{3.20}$$

and

$$\Gamma_{\chi_0\to V+\gamma} = 4\alpha M_\chi\left(\frac{m_\chi^2 - m_V^2}{2m_\chi^2}\right)g^2_{\chi_0V} \tag{3.21}$$

where m_χ is the mass of pseudoscalar χ_0 meson.

4. RELATIVISTIC EQUATION FOR BARYONS [1]

We apply the above developed method to obtain an equation describing the baryons. In the quark model the baryons are regarded as bound states of three quarks. Thus, the wave function of the baryon is a spinor of the third rank $\psi_{ABC}(x_1, x_2, x_3)$.

By analogy with the meson case we consider the relativistic invariant equation for ψ_{ABC} which admits the solution of the form

$$\psi_{ABC}(x_1, x_2, x_3) = \varphi(x_1 - x_3, x_2 - x_3)\phi_{ABC}(x_1 + x_2 + x_3), \tag{4.1}$$

where ϕ_{ABC} is the spinor of the third rank which describes the motion of the centre of mass of the bound system of three quarks, i.e., the baryon as a whole, φ stands for some scalar function.

To get an equation describing the motion of baryons in a weak external electromagnetic field we shall start from the squared Dirac equation

$$\begin{aligned} \mathcal{D}_A(x_1)\mathcal{D}_B(x_2)\mathcal{D}_C(x_3)\psi_{ABC}(x_1, x_2, x_3) = g_0 W(x_1, x_2, x_3) \times \\ \times \int dx_1' dx_2' dx_3' W(x_1', x_2', x_3')\delta(x_1 + x_2 + x_3 - x_1' - x_2' - x_3')\psi_{ABC}(x_1', x_2', x_3'), \end{aligned} \tag{4.2}$$

where the operators $\mathcal{D}$ are given by the expression (3. 3), $W(x_1, x_2, x_3)$ is the scalar function which satisfies the conditions

$$\begin{aligned} W(x_1, x_2, x_3) &= W(x_1 + a, x_2 + a, x_3 + a), \\ W(x_1, x_2, x_3) &= W(-x_1, -x_2, -x_3). \end{aligned} \tag{4.3}$$

The Fourier transform $W(q_1, q_2, q_3)$ of the scalar function W may be, for example, taken in the form

$$W(q_1, q_2, q_3) = \sum_N \frac{C_N}{[N^2 - (q_1 - q_2)^2 - i\epsilon][N^2 - (q_2 - q_3)^2]} \times \frac{1}{N^2 - (q_1 - q_3)^2 - i\epsilon}. \tag{4.4}$$

To get the equation for the function

$$\begin{aligned} \phi_{ABC}(x_1 + x_2 + x_3) &= \int W(x_1', x_2', x_3')\psi_{ABC}(x_1', x_2', x_3') \times \\ &\times \delta(x_1 + x_2 + x_3 - x_1' - x_2' - x_3')dx_1' dx_2' dx_3' \end{aligned} \tag{4.5}$$

we use again, by analogy with the meson case, the expansion in the large mass of the quark M and retain the terms of the first order in the charge e. As a result, for the Fourier transform of the function (4. 5) we get the following equation

$$\begin{aligned} (p^2 - m_N^2)\phi_{ABC}(p^2) &= 2\int dk\, f(k^2) A \cdot p[e_A + e_B + e_C]\phi_{ABC}(p - k) \times \\ &\times 3\int dk\, f(k^2)[(e\hat{k}\hat{A})_A + (e\hat{k}\hat{A})_B + (e\hat{k}\hat{A})_C]\phi_{ABC}(p - k), \end{aligned} \tag{4.6}$$

where the nucleon mass m_N is determined from the equation similar to (2.13), while the function $f(k^2)$ is given by

$$f(k^2) = \frac{\int W(p_1,p_2,p_3)W(p_1+k,p_2,p_3)\delta(p_1+p_2+p_3)dp_1\,dp_2\,dp_3}{\int W^2(p_1,p_2,p_3)\delta(p_1+p_2+p_3)dp_1\,dp_2\,dp_3}. \tag{4.7}$$

From this, in particular, it is seen that* $F(0) = 1$.

When the electromagnetic field is absent Eq. (4.2) is invariant under $\tilde{U}(12)$ as the operators $\mathcal{D}$ and W have no spin and unitary structure; the free equation

$$(p^2 - m^2)\phi_{ABC} = 0 \tag{4.8}$$

is also invariant under $\tilde{U}(12)$. In order not to mix the solutions with positive and negative energies in the free case, we use a supplementary condition

$$(\hat{p})_A^{A'}\phi_{A'BC} = (\hat{p})_B^{B'}\phi_{A'B'C} = (\hat{p})_C^{C'}\phi_{ABC'}. \tag{4.9}$$

We require that ϕ_{ABC} satisfy simultaneously the following equations

$$(\hat{p}-m)_A\phi_{ABC} = (\hat{p}-m)_B\phi_{ABC} = (\hat{p}-m)_C\phi_{ABC} = 0. \tag{4.10}$$

Conditions (4.9) and (4.10) break the invariance under the group $\tilde{U}(12)$, conserving, at the same time, the invariance under $SU(6)$.

Consider the solution (4.8)–(4.10) corresponding to the symmetrical spinor of the third rank. The symmetrical spinor ϕ_{ABC} is a relativistic generalization of 56-plet in the non-relativistic scheme $SU(6)$.

The symmetrical spinor ϕ_{ABC} may be represented as [2]

$$\begin{aligned}\phi_{ABC} = \phi_{(\alpha p)(\beta q)(\gamma r)} &= \sqrt{\frac{1}{8}}D_{\alpha\beta\gamma}d_{pqr} + \\ &+ \frac{1}{6\sqrt{2}}\left[N_{[\alpha\beta]\gamma}\epsilon_{pqs}B_r^s + N_{[\beta\gamma]\alpha}\epsilon_{qrs}B_p^s + N_{[\gamma\alpha]\beta}\epsilon_{rps}B_q^s\right],\end{aligned} \tag{4.11}$$

where $D_{\alpha\beta\gamma}$ is the symmetrical spinor of the third rank, $N_{[\alpha\beta]\gamma}$ is the spinor antisymmetrical with respect to the indices α, β, d_{pqr} stands for the unitary decuplet, B_q^p is the unitary octet.

Using relations (2.17) one can obtain [2]

$$D_{\alpha\beta\gamma} = \psi_{\alpha\mu}(\gamma_\mu C)_{\beta\gamma} - \frac{i}{2m}(p_\mu\psi_\nu - p_\nu\psi_\mu)_\alpha(\sigma^{\mu\nu}C)_{\beta\gamma}, \tag{4.12a}$$

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*To get form factors of mesons and baryons (3.8–4.6) we use the equation which contains the squared Dirac operators $\hat{\partial}^2$, and the factorizable potentials. As Bogolubov has shown [3], these results remain unchanged, if we use nonsquared Dirac operators for quarks and describe bound states with a Bethe–Salpeter-type equation. For example Bogolubov's equation for bosons has the form

$$\left\{(\hat{\partial}_1 - M)(\hat{\partial}_2 + M) - \gamma_1^{(5)}\gamma_2^{(5)}W\right\}\psi = 0,$$

where

$$W(x_1,x_2) = M^2 - u(x_1 - x_2).$$

$$N_{[\alpha\beta]\gamma} = \frac{[(\hat{p}+m)\gamma_s C]_{\alpha\beta}}{m}\,\psi_\gamma, \tag{4.12b}$$

where C is determined by $(\gamma_\mu C)_{\alpha\beta} = (\gamma_\mu C)_{\beta\alpha}$, $C^T = -C$.

The spinors ψ and ψ_μ satisfy the equations

$$(\hat{p}-m)\psi_\mu = 0, \quad \gamma^\mu\psi_\mu = 0, \tag{4.13a}$$

$$(\hat{p}-m)\psi = 0, \quad p^\mu\psi_\mu = 0, \tag{4.13b}$$

so:

$$D_{\alpha\beta\gamma} = \psi_{\mu\alpha}\left[\left(1+\frac{\hat{p}}{m}\right)\gamma_\mu C\right]_{\beta\gamma}.$$

Consider the baryon in a constant magnetic field. Retaining only the large components of the spinor ϕ_{ABC} we get from (4.6)

$$(p^2-m^2)\phi_{ABC} = 3[(e\vec{\sigma})_A + (e\vec{\sigma})_B + (e\vec{\sigma})_C]\vec{H}\phi_{ABC}. \tag{4.14}$$

Here the spinor indices run from one to two. It can be seen from (4.14) that the effective magnetic moment of the quark with charge e in the baryon is $3e/2m$ and the operator of the baryon magnetic moment is

$$\vec{\mu} = \frac{3}{2m}[(e\vec{\sigma})_1 + (e\vec{\sigma})_2 + (e\vec{\sigma})_3]. \tag{4.15}$$

Averaging (4.14) over the functions corresponding to the particles with spins 3/2 and 1/2 we get the well-known relations among the magnetic moments following from the $SU(6)$ symmetry [4].

It is rather interesting that the magnetic moment of the proton calculated by Eq.(4.14) turns out to be equal to three nuclear magnetons. Thus, just as in the meson case, the magnetic moment in the composite model is determined by the mass of the bound state, and the composite particle possesses an anomalous magnetic moment.

Let us now study the electromagnetic form factors of baryons. We consider the vertex

$$3\bar{\phi}^{ABC}\Gamma_A^{A'}\phi_{A'BC} = 2mJ_\alpha A^\alpha, \tag{4.16}$$

where

$$\Gamma_A^{A'}(k) = \left[2e_p^{p'}A_\mu\cdot p^\mu + \frac{3}{2}(\sigma^{\mu\nu})_\alpha^{\alpha'}e_p^{p'}F_{\mu\nu}(k)\right]f(k^2)$$

is the matrix of the electric charge.

Substituting (4.11), (4.12) into (4.16) we obtain

$$\begin{aligned} J_\alpha &= 3\left\{-\frac{p_\alpha}{m}\left(1+\frac{k^2}{4m^2}\right)(\bar{\psi}_\mu\psi^\mu) + \frac{p_\alpha}{m^3}\left(k^\mu k^\nu - \frac{1}{4}g^{\mu\nu}k^2\right)\bar{\psi}_\mu\psi_\nu + \right.\\ &+ \left.\frac{3kp}{2m}(\bar{\psi}_\mu\psi_\alpha - \bar{\psi}_\alpha\psi_\mu)\right\} d^{pqr}e_p^{p'}d_{p'qr}f(k^2)\times\\ &\times \left\{\left(1+\frac{k^2}{4m^2}\right)\frac{p_\alpha}{m}(\bar{\psi}\psi)_F^e - \frac{1}{4m^2}(\bar{\psi}r_\alpha\psi)_{3D+2F}^e\right\} f(k^2) + \\ &+ \frac{3}{m^2}\epsilon^{\alpha\beta\sigma\rho}p_\sigma k_\rho\left\{\bar{\psi}_\beta\psi d^{pqr}e_p^{p'}\epsilon_{p'qs}B_r^s - (\bar{\psi}\psi_\beta)\epsilon^{pqs}B_s^r e_p^{p'}d_{p'qr}\right\} f(k^2) \end{aligned} \tag{4.17}$$

$$r_\alpha = 2\epsilon^{\alpha\sigma\rho\lambda} p_\sigma k_\rho \gamma_\lambda \gamma_5.$$

Here d_{pqr} is the wave function of the decuplet, B_r^s is the wave function of the octet, ψ_μ is the wave function of particles with spin 3/2, and ψ is the wave function for particles with spin 1/2. From (4.17) we have

$$\begin{aligned}
G_E^{3/2} &= eF(k^2)\left(1+\frac{k^2}{4m^2}\right), \\
G_E^{1/2} &= eF(k^2)\left(1+\frac{k^2}{2m^2}\right), \\
G_M^{3/2} &= \frac{3e}{2m}F(k^2), \\
G_M^{1/2} &= \frac{e}{2m}F(k^2)\mu; \quad \mu_p = 3 \quad \mu_n = -2 \quad \mu_{\Sigma^-} = -1 \\
G_Q^{2/2} &= \frac{e}{m^3},
\end{aligned} \tag{4.18}$$

where $G_E^{3/2,1/2}$ are the electric form factors of baryons with spins 3/2 and 1/2, respectively $G_M^{3/2,1/2}$ are the magnetic form factors, $G_Q^{3/2}$ is the quadrupole moment.

The last term describes the photoproduction and decay of resonances. The matrix element for the radiative decay is given by the expression

$$M = \frac{6e}{m} g_{d\to b+\gamma} \epsilon^{\alpha\beta\sigma\rho} p_\sigma k_\rho A_\alpha \bar{\psi}\psi_\beta. \tag{4.19}$$

For the constants g_{db} we have [5]

$$g_{\Delta^+ p} = g_{\Delta^0 n} = g_{\Xi^0\Xi^0} = g_{\Sigma^+\Sigma^+} = -\frac{2}{\sqrt{3}} g_{\Sigma^0\Lambda} = 2g_{\Sigma^0\Sigma^0} = 1. \tag{4.20}$$

Taking into account the mass difference only in the phase space, we get for the decay width of resonances

$$\Gamma_{d\to b+\gamma} = 3\alpha\frac{m_d^2}{m_b}\left(\frac{m_d^2+m_b^2}{2m_d^2}\right)^3\left(1+\frac{m_b}{m_d}\right)^2, \tag{4.21}$$

where m_d is the mass of the baryon from the decuplet, and m_b is the mass of the baryon from the octet. We have seen that the interaction of the composite particle with an external electromagnetic field is nonlocal.

Note that if we apply these results to the examination of vertices in quantum electrodynamics we have to restrict ourselves to the first order of the perturbation theory, in order not to deal with difficulties connected with violation of unitarity and causality due to the nonlocal character of the interaction.

Now we discuss in more detail the symmetry properties of the wave functions of baryons. The equation which describes the baryons admits a solution of the form

$$\psi_{ABC}(x,y,z) = \varphi(x,y,z)\phi_{ABC}(x+y+z), \tag{4.22}$$

where φ is a scalar function which is symmetric with respect to the variables x, y, z.

In studying the form factors we have used the wave function which is symmetric with respect to the permutation of indices A, B, C. So we choose the completely symmetric wave function ψ_{ABC} to describe the baryon state.

However, if we consider the quarks to be Dirac particles which can be also in free states, we must require that the wave functions of baryons be antisymmetric with respect to the permutation of all quantum numbers of true quarks. Here we meet a contradiction. Choosing the completely symmetric wave function for the baryons we get good agreement between the theoretical and experimental results, although such a choice contradicts the Pauli principle. In order to avoid this contradiction we assign a new quantum number to the quark*. We assume that, besides the spin and unitary index, the quark is characterized by the quantum number a and the quark wave function will be denoted by $\psi_{\underline{ABC}}(x,y,z) = \psi_{ABC;abc}(x,y,z)$ where a runs from one to three. For the baryon wave function we have

$$\psi_{\underline{ABC}}(x,y,z) = \psi_{ABC,abc}(x,y,z)$$

We determined that the equation which describes the baryon allows a solution of the form

$$\psi_{\underline{ABC}}(x,y,z) = \frac{1}{\sqrt{6}}\epsilon_{abc}\varphi(x,y,z)\psi_{ABC}(x,y,z), \tag{4.23}$$

where ϵ_{abc} is the totally antisymmetric tensor. The presence of the tensor ϵ_{abc} guarantees the antisymmetric property of the baryon functions, if the function $\psi_{ABC}(x,y,z)$ is completely symmetric.

The relativistic invariant equation which allows the solution of the form (4.23) will be obtained by a strict generalization of Eq. (4.2). It has the form

$$\mathcal{D}_{\underline{A}}^{\underline{A'}}\mathcal{D}_{\underline{B}}^{\underline{B'}}\mathcal{D}_{\underline{C}}^{\underline{C'}}\psi_{\underline{A'}\underline{B'}\underline{C'}}(x,y,z) = \tfrac{1}{6}g\epsilon_{abc}W(x,y,z) \times$$
$$\times \int dx'dy'dz'W(x',y',z')\delta(x+y+z-x'-y'-z')\epsilon^{a'b'c'}\psi_{ABC,a'b'c'}(x',y',z'), \tag{4.24}$$

where D is the squared Dirac operator. In order to describe the baryons in an electromagnetic field, we must make a usual substitution

$$i\hat{\partial}_x \to i\hat{\partial}_x - e\hat{A}(x), \tag{4.25}$$

where e is the charge operator. We suppose that the charge operator acts not only on the unitary index but it acts also on the additional quantum number. Such a generalization of the charge operator leads to an integer value for the quark charge. We want to emphasize that the prediction of Eq.(4.2) concerning the form factors remains unchanged if we treat three triplets on the same footing.

We define the charge operator in the following form

$$e_{\underline{A}}^{\underline{A'}} = e_A^{A'} + e_a^{a'}. \tag{4.26}$$

The operator $e_A^{A'}$ acts on the unitary index and has the form

$$e_A^{A'} = \begin{pmatrix} \frac{2}{3} & 0 & 0 \\ 0 & -\frac{1}{3} & 0 \\ 0 & 0 & -\frac{2}{3} \end{pmatrix}. \tag{4.27}$$

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*This possibility is also discussed in the recent paper by Nambu [4].

The operator $e_a^{a'}$ acts on the additional quantum numbers and is of the form

$$e_a^{a'} = \begin{pmatrix} \frac{1}{3} & 0 & 0 \\ 0 & \frac{1}{3} & 0 \\ 0 & 0 & -\frac{2}{3} \end{pmatrix}. \tag{4.28}$$

Then the charge of the quark is integer.

$$\begin{array}{c|ccc} {}_{A}\,{}^{a} & 1 & 2 & 3 \\ \hline 1 & 1 & 1 & 0 \\ 2 & 0 & 0 & -1 \\ 3 & 0 & 0 & -1 \end{array}. \tag{4.29}$$

Taking into account that the operator $e_a^{a'}$ is diagonal and Sp $e_a^{a'} = 0$ we get

$$\epsilon^{abc} e_a^{a'} \epsilon_{a'bc} = 0. \tag{4.30}$$

Therefore, we have the identity

$$e_A^{A'} = \frac{1}{6} \epsilon^{abc} \left(e_A^{A'} + e_a^{a'} \right) \epsilon_{a'bc}. \tag{4.31}$$

Let us consider Eq.(4.26) where the charge operator e has the form $e_{\underline{A}}^{\underline{A}'} = e_A^{A'} + e_a^{a'}$. Retaining in this equation the first order terms in e and making use of the expansion in the large mass of the quark M, we get an equation describing the motion of the baryon in a weak electromagnetic field

$$\begin{aligned} (p^2 - m^2)\phi_{\underline{ABC}}(p) &= \int dk\, f(k) \left[\Gamma_{\underline{A}}^{\underline{A}'} + \Gamma_{\underline{B}}^{\underline{B}'} + \Gamma_{\underline{C}}^{\underline{C}'} \right] \phi_{\underline{A}'\underline{B}'\underline{C}'}(p-k) \\ \Gamma_{\underline{A}}^{\underline{A}'} &= [2A \cdot pe + 3(e\hat{k} \cdot \hat{A})]_{\underline{A}}^{\underline{A}'}. \end{aligned} \tag{4.32}$$

For the nonperturbed solution we have

$$\phi_{\underline{ABC}}^{(p)} = \frac{1}{\sqrt{6}} \epsilon_{abc} \phi_{ABC}(p). \tag{4.33}$$

Substituting (4.33) into (4.32) after averaging over the functions $\frac{1}{\sqrt{6}}\epsilon_{abc}$ and taking into account (4.32) we arrive again to Eq.(1.2) which was our goal.

Assigning a new quantum number, we may construct 3 types of mesons $\psi_{A,1}^{B,1}$, $\psi_{A,2}^{B,2}$, $\psi_{A,3}^{B,3}$. The charges of these mesons coincide with the charge of the known mesons. The mesons are constructed by the triplet 1 with charge (100), the mesons $\psi_{A,2}^{B,2}$ — for the triplet 2 with charge (100) and $\psi_{A,3}^{B,3}$ are constructed by the triplet 3 with charge (0 -1 -1).

In order that the meson form factors remained unchanged, it is sufficient that the zero approximation solution take the form

$$\psi_{A,a}^{B,b} = \frac{1}{\sqrt{3}} \delta_b^a \psi_A^B. \tag{4.34}$$

Such a solution selects the mesons in which the states ψ_1^1, ψ_2^2, ψ_3^3 appear with equal weights.

Now I want to discuss briefly a possible test of the hypothesis that the mesons are constructed with three triplets.

From the expression for the electromagnetic meson form factors it is easy to see that assigning the new quantum number we change only the radiative decay vertex.

Indeed, taking the charge operator in the form $\hat{e} \to \hat{e} + e_a$ we get

$$\begin{aligned} [\bar{\varphi}_5 \varphi_\beta]_F^e &\to [\bar{\varphi}_5 \varphi_\beta]_F^e \\ [\bar{\varphi}_5 \varphi_\beta]_D^e &\to [\bar{\varphi}_5 \varphi_\beta]_D^e + 2e_a \bar{\varphi}_5 \varphi_\beta, \end{aligned} \tag{4.35}$$

where $e_a = 1/3$ for the mesons ψ_1^1, ψ_2^2, and $-2/3$ for the mesons ψ_3^3.

The results of calculations for the radiative decay constant are given in the Table.

Table. RADIATIVE DECAY CONSTANTS

Decay	Quark models	ψ_1^1, ψ_2^2	ψ_3^3
$\rho_0 \to \pi_0 \gamma$	$1/3$	1	-1
$k^{*\pm} \to k^\pm \gamma$	$1/3$	1	-1
$k^{*0} \to k^0 \gamma$	$-2/3$	0	-2
$\omega \to \pi^0 \gamma$	1	1	2
$\omega \to \eta^0 \gamma$	$\frac{1}{3\sqrt{3}}$	$\frac{1}{\sqrt{3}}$	$-\frac{1}{\sqrt{3}}$
$\varphi \to \eta_0 \gamma$	$\frac{2}{3}\sqrt{\frac{2}{3}}$	0	$2\sqrt{\frac{2}{3}}$
$\chi_0 \to \rho_0 \gamma$	$-\sqrt{\frac{2}{3}}$	$-\sqrt{\frac{2}{3}}$	$-\sqrt{\frac{2}{3}}$
$\chi_0 \to \omega \gamma$	$-\frac{2}{3}\sqrt{\frac{2}{3}}$	0	$-2\sqrt{\frac{2}{3}}$

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КРАТКИЙ ОЧЕРК ЖИЗНИ И ДЕЯТЕЛЬНОСТИ АКАДЕМИКА А.Н. ТАВХЕЛИДЗЕ

Альберт Никифорович Тавхелидзе родился 16 декабря 1930 г. в Тбилиси. Он рос в дружной гостеприимной семье с богатыми традициями, очень музыкальной, глубоко почитавшей людей культурных и образованных.

В 1948 г. А.Н. Тавхелидзе окончил 8-ю мужскую среднюю школу г.Тбилиси, а в 1953 г. Тбилисский государственный университет по специальности «теоретическая физика». Полученная в университете основательная подготовка и рекомендации академиков И.Н. Векуа и Н.И. Мусхелишвили дали ему возможность пройти обучение в аспирантуре в Математическом институте им. В.А. Стеклова Академии наук СССР. Его научным руководителем был академик Н.Н. Боголюбов. В 1957 г. на Ученом совете МИАН А.Н. Тавхелидзе успешно защитил кандидатскую диссертацию на соискание ученой степени кандидата физико-математических наук.

С 1956 г., по приглашению Н.Н. Боголюбова и А.А. Логунова, А.Н. Тавхелидзе начал свою научную деятельность в Лаборатории теоретической физики Объединенного института ядерных исследований (ЛТФ ОИЯИ), только что созданного в Дубне международного физического центра. Высокая научная, исключительно благоприятная и доброжелательная атмосфера, царившая в ОИЯИ, способствовала творческому научному поиску.

В Дубне А.Н. Тавхелидзе работал в период 1956–1970 гг. и прошел путь от научного сотрудника до заместителя директора Лаборатории теоретической физики. В 1963 г. А.Н. Тавхелидзе на Ученом совете ЛТФ ОИЯИ защитил диссертацию на соискание ученой степени доктора физико-математических наук. В 1965 г. ему было присвоено звание профессора. Вместе с академиками Н.Н. Боголюбовым и А.А. Логуновым он много внимания уделял подготовке научной смены, привлекая наиболее одаренных студентов и аспирантов в Лабораторию теоретической физики.

А.Н. Тавхелидзе принадлежит большая заслуга в организации и проведении в Дубне и за ее пределами целого ряда симпозиумов, школ, конференций и пр. Особо следует отметить неформальные семинары, регулярно проводившиеся в 1966–1971 гг. ОИЯИ и ЦЕРН (Европейским центром ядерных исследований в Женеве). Целью этих встреч было обсуждение наиболее перспективных направлений исследований в области физики высоких энергий. В работе одного из таких семинаров, проходившего в Тбилиси в 1969 г., принимали участие директора почти всех ядерных центров мира.

В период работы в ЛТФ А.Н. Тавхелидзе активно участвует в научной и научно-организационной работе других физических институтов.

В 1965–1970 гг. А.Н. Тавхелидзе по приглашению академика А.А. Логунова принял участие в организации сектора теоретической физики Института физики высоких энергий (Протвино) и был его первым руководителем. В результате совместных усилий Н.Н. Боголюбова и А.Н. Тавхелидзе в Киеве в 1967 г. был открыт Институт теоретической физики (ныне носящий имя Н.Н. Боголюбова), где А.Н. Тавхелидзе работал заведующим отделом физики элементарных частиц (1967–1971 гг.).

В 1967 г., в тесном сотрудничестве с академиками Н.Н. Боголюбовым и М.А. Марковым, он добивается правительственного решения об издании журнала АН СССР «Теоретическая и математическая физика». До 1991 г. А.Н. Тавхелидзе являлся заместителем главного редактора этого журнала.

В 1970 г. президиум Академии наук СССР утвердил А.Н. Тавхелидзе в должности директора вновь созданного Института ядерных исследований (Москва), поручив ему разработку структуры и формирование научной тематики этого центра. После дискуссий с участием видных ученых-экспертов, при решающей поддержке академика М.А. Маркова, в ИЯИ АН СССР сформировались два основных научных направления: физика частиц и атомного ядра и нейтринная астрофизика. Было запланировано создание базовых ядерно-физических установок Института: Московской мезонной фабрики (в г. Троицке) и нейтринных обсерваторий с соответствующими нейтринными телескопами, подземными — в Приэльбрусье (Баксанская нейтринная обсерватория) и подводным на озере Байкал. Фактически тематика Института представляла собой одно новое направление — частицы и космология.

К 1986 г., когда А.Н. Тавхелидзе был избран президентом Академии наук Грузии, в Институте была почти завершена вся предварительная работа, развернуто широкое международное сотрудничество с научными центрами Италии, Америки и других стран, начаты первые эксперименты. Теоретические и экспериментальные исследования сотрудников этого Института были уже хорошо известны широкой научной общественности. В Институте ядерных исследований А.Н. Тавхелидзе продолжает плодотворно работать по сей день, являясь заведующим отделом теоретической физики — научным руководителем ИЯИ.

Начиная с 1969 г., раз в два года, в основном на базе ИЯИ РАН и АН Грузии, А.Н. Тавхелидзе проводит международные конференции, посвященные физике кварков.

В 1970–1986 гг. А.Н. Тавхелидзе — профессор Московского государственного университета им. М.В. Ломоносова, что, несомненно, сыграло важную роль при формировании широко известного в научном мире отдела теоретической физики ИЯИ РАН.

В период работы в Москве активизировались связи А.Н. Тавхелидзе с Тбилисским государственным университетом и Академией наук Грузии. В 1971 г. при поддержке академика Н.И. Мусхелишвили в Математическом институте им. А.М. Размадзе был создан отдел теоретической физики, целиком укомплектованный молодыми учеными, прошедшими стажировку в Дубне. Руководство отделом было поручено А.Н. Тавхелидзе.

В 1980 г. на базе проблемной лаборатории физики высоких энергий, существовавшей при физическом факультете Тбилисского университета (ТГУ), где работали в основном физики, выросшие в Дубне и поддерживающие с ней научные связи, был создан Институт физики высоких энергий ТГУ. Научное руководство институтом было возложено на А.Н. Тавхелидзе.

Вклад грузинских ученых в развитие физики высоких энергий был по достоинству оценен, и в 1976 г., при большой поддержке правительства Грузии, в Тбилиси состоялась XVIII Рочестерская конференция по физике высоких энергий, в которой участвовало почти 1500 физиков.

В 1967 г. А.Н. Тавхелидзе был избран членом-корреспондентом Академии наук Грузии, а в 1974 г. — ее действительным членом.

В 1986 г., по рекомендации правительства Грузии и при поддержке Академии наук СССР, А.Н. Тавхелидзе был избран президентом Академии наук Грузии. В 1993 и 1998 гг. Общее собрание Академии наук Грузии вновь подтвердило его полномочия как президента. Вклад Академии наук в научную и интеллектуальную жизнь Грузии получил должную оценку. По представлению президента страны парламент Грузии в июне 1999 г. принял Закон об Академии наук Грузии. Законом установлен государственный статус Академии наук Грузии, определены ее задачи и гарантии правовой, имущественной, организационной и финансовой деятельности.

С 1986 г. А.Н. Тавхелидзе является главным редактором журнала «Сообщения Академии наук Грузии», который в настоящее время издается на грузинском и английском языках.

С 1995 г. он директор Института физики высоких энергий ТГУ им. И.А.Джавахишвили. Накопленные в институте знания и опыт сделали возможным формирование сети Интернет в АН Грузии и государственных высших учебных заведениях Грузии.

А.Н. Тавхелидзе — автор более двухсот научных публикаций, которые характеризуются высоким индексом цитируемости. Многие из его соавторов были начинающими научными сотрудниками, а сегодня они — признанные в мире крупные ученые, организаторы науки.

За выдающийся вклад в науку А.Н. Тавхелидзе в 1984 г. был избран членом-корреспондентом Академии наук СССР, а в 1990 г. — действительным членом (академиком) АН СССР (с 1991 г. — Российской академии наук).

А.Н. Тавхелидзе — лауреат Государственной премии СССР (1973 г.), которой был отмечен цикл совместных работ «Фоторождение π-мезонов на нуклонах». Цикл совместных исследований «Новое квантовое число — цвет и установление динамических закономерностей в кварковой структуре элементарных частиц и атомных ядер» в 1988 г. был удостоен Ленинской премии. В 1998 г. «за создание Баксанской нейтринной обсерватории и исследования в области нейтринной астрофизики элементарных частиц и космических лучей» А.Н. Тавхелидзе, совместно с сотрудниками Института ядерных исследований, была присуждена Государственная премия Российской Федерации.

В 1987 г. в Государственном реестре открытий СССР было зарегистрировано открытие «Правила кваркового счета Матвеева–Мурадяна–Тавхелидзе».

Национальная академия наук Украины в 1996 г. присудила А.Н. Тавхелидзе премию имени Н.Н. Боголюбова. Международная ассоциация академий наук, принимая во внимание большой вклад А.Н. Тавхелидзе в укрепление международного научного сотрудничества, в 1998 г. наградила его золотой медалью. За активное участие в работе Всемирной федерации ученых и в связи с празднованием 2000-летия Рождества Христова Папа Римский Иоанн Павел II вручил А.Н. Тавхелидзе памятный знак.

В 1987–1990 гг. А.Н. Тавхелидзе был депутатом Верховного Совета Грузии и членом Президиума Верховного Совета. В 1989 г. он был избран народным депутатом СССР.

А.Н. Тавхелидзе имеет ряд высших государственных наград СССР и Российской Федерации. Он член нескольких иностранных академий, участник Пагуошского движения ученых за мир.

Сам Альберт Никифорович в автобиографии так характеризует свой путь в науке: «Как физик я вырос в школе академика Н.Н. Боголюбова — моего наставника и учителя. Пройдя большой путь научной и научно-организационной деятельности, тесно сотрудничал и общался с выдающимися учеными, государственными и хозяйственными деятелями крупного масштаба. Большинство из них — неповторимые личности, люди с огромным жизненным опытом, встречи и общение с которыми были весьма впечатляющими и оказали большое влияние на формирование моей личности».

А.Н. Тавхелидзе и его супруга Майа Меунаргия имеют дочь Нато и сына Ношревана, а также внуков — Майко и Нику...

Все, кто хорошо знает Альберта Никифоровича, всегда отмечают такие его качества, как преданность идеалам науки, огромную силу воли, редкую работоспособность и умение сплотить людей для достижения поставленной цели, надежность в дружбе и доброту.

ТРУДЫ АКАДЕМИКА А.Н. ТАВХЕЛИДЗЕ ПО ОСНОВНЫМ НАПРАВЛЕНИЯМ ЕГО НАУЧНОЙ ДЕЯТЕЛЬНОСТИ
(с краткими комментариями)

1. Дисперсионные соотношения (ДС) и приближенные уравнения в квантовой теории поля (КТП)

Обобщая метод ДС Н.Боголюбова для неупругих процессов и процессов с переменным числом частиц в КТП, А.Тавхелидзе совместно с А.Логуновым впервые получил ДС для амплитуд фоторождения π-мезонов на нуклонах. На основе данных ДС в приближении двухчастичной унитарности были получены сингулярные интегральные уравнения, типа уравнений Н.Мусхелишвили, для парциальных амплитуд фоторождения (совместно с А.Логуновым и Л.Соловьевым). Ядром этих уравнений является амплитуда мезон-нуклонного рассеяния, а неоднородная часть представляет собой значение амплитуды фоторождения в однонуклонном приближении. Эти результаты получили экспериментальное подтверждение.

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2. Квазипотенциальный метод в КТП

А.Тавхелидзе совместно с А.Логуновым предложил трехмерную формулировку КТП, в рамках которой для описания системы двух взаимодействующих частиц получены релятивистские трехмерные квазипотенциальные уравнения, известные в литературе как уравнения Логунова–Тавхелидзе. Был развит регулярный метод построения квазипотенциала (КП) на основе двухчастичной функции Грина. Благодаря вкладу многочастичных промежуточных состояний квазипотенциал является комплексной функцией энергии. Установлены аналитические свойства КП по энергии.

Уравнения Логунова–Тавхелидзе с успехом применяются для вычисления сверхтонких поправок к энергии атома водорода, энергии связанных состояний кварков и т.п. Многочастичные квазипотенциальные уравнения, записанные в переменных светового фронта, удобно использовать для анализа кварковой структуры адронов и ядер, изучения процессов при больших передачах импульса.

В работах, выполненных совместно с В.Гарсеванишвили, А.Квинихидзе, В.Матвеевым, А.Сисакяном и Л.Слепченко, на базе описания релятивистских составных систем в переменных светового фронта изучены процессы множественного образования частиц и впервые строго установлен асимптотический закон P_T^{-4} для инклюзивного образования адронов с большими поперечными импульсами.

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3. Конечноэнергетические (КЭ) правила сумм и дуальность

А.Тавхелидзе совместно с А.Логуновым и Л.Соловьевым получены КЭ правила сумм для амплитуды мезон-нуклонного рассеяния и на их основе установлено свойство глобальной дуальности — интегральные соотношения между резонансной частью амплитуды рассеяния и реджевскими параметрами. Свойства глобальной и локальной (Венециано) дуальности послужили основой для формулировки струнной модели адронов. В дальнейшем совместно с Н.Красниковым и К.Четыркиным метод КЭ правил сумм был обобщен на случай квантовой хромодинамики с использованием характерного для этой теории свойства асимптотической свободы.

КЭ правила сумм являются непертурбативным методом, широко используемым для вычислений в квантовой хромодинамике.

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4. Массы фермионов и явление спонтанного нарушения симметрии

Микроскопическая теория сверхтекучести и сверхпроводимости Н.Боголюбова основана на открытом им явлении спонтанного нарушения симметрии основного состояния. А.Тавхелидзе одним из первых (совместно с Б.Арбузовым и Р.Фаустовым) установил возможность возникновения массы фермионов за счет спонтанного нарушения симметрии в двумерной модели КТП.

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5. Квантовое число цвет. Физическая модель адронов как связанных состояний цветных кварков

В 1965 г. А.Тавхелидзе совместно с Н.Боголюбовым и Б.Струминским, независимо от Намбу и Хана, выдвинул гипотезу о наличии у кварков нового квантового числа, названного впоследствии «цветом». Согласно этой гипотезе каждый тип кварков с данным ароматом может находиться в трех унитарно эквивалентных состояниях, соответствующих трем значениям цвета. Одновременно был сформулирован принцип отбора физических состояний квантовых систем, описывающих наблюдаемые мезоны и барионы, который отвечает требованию бесцветности адронных состояний. Было отмечено также, что цветные кварки могут иметь как дробные, так и целочисленные электрические заряды. В последнем случае цветовая симметрия должна была бы нарушаться по крайней мере в электромагнитных (ЭМ) взаимодействиях.

В отличие от модели кварков Гелл-Мана и Цвейга в динамическом подходе Боголюбова–Тавхелидзе цветные кварки рассматривались как физические фундаментальные частицы — фермионы, находящиеся в адронах в квазисвободном состоянии. Вычисленные в рамках этого подхода ЭМ- и слабые формфакторы адронов, ширины ЭМ-распадов мезонов получили экспериментальное подтверждение, что оправдало заложенную в основу динамической модели гипотезу квазисвободных кварков.

В рамках модели квазинезависимых кварков были получены правила «кваркового счета Матвеева–Мурадяна–Тавхелидзе», согласно которым при больших энергиях и передачах импульса имеет место степенное поведение по большому импульсу амплитуды упругого рассеяния и формфакторов адронов. При этом показатель степени выражается через полное число кварков, из которых составлены частицы, участвующие в реакциях. Обобщение формул кваркового счета на другие процессы с участием адронов и лептонов и их экспериментальная проверка является существенным этапом в понимании физической природы цветных кварков и кварковой структуры адронов.

Динамическая кварковая модель адронов легла в основу поисков релятивистского обобщения $SU(6)$-симметрии элементарных частиц и привела к формулировке релятивистски-ковариантных уравнений для связанных систем частиц в квантовой теории поля (совместные работы с В.Кадышевским, Р.Мурадяном, Нгуен Ван Хьеу, И.Тодоровым).

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6. Масштабная инвариантность процессов при высоких энергиях. Принцип автомодельности

В цикле работ А.Тавхелидзе, выполненных совместно с В.Матвеевым и Р.Мурадяном, выдвинут принцип автомодельности в физике высоких энергий (1969 г.) и на его основе развит единый подход к описанию явлений масштабно-инвариантного поведения различных процессов глубоконеупругого взаимодействия лептонов с адронами.

Предсказано масштабно-инвариантное поведение процессов образования пар мюонов в столкновениях адронов при высоких энергиях в области больших инвариантных масс димюонов (названных позднее процессами Дрелла–Яна).

Дано обобщение принципа автомодельности для процессов сильного взаимодействия адронов при высоких энергиях, в частности, для процессов инклюзивного образования адронов с большими поперечными импульсами.

В совместных работах с Н.Боголюбовым и В.Владимировым (1972 г.) в рамках локальной квантовой теории поля дано строгое обоснование существования автомодельных (масштабно-инвариантных) асимптотик глубоконеупругих процессов и установлена точная взаимосвязь структурных функций этих процессов с поведением коммутаторов локальных токов в окрестности светового конуса.

В последующих работах (совместно с Э.Вицореком, В.Матвеевым и Д.Робашеком) дано обобщение этих результатов для недиагональных матричных элементов T-произведений локальных токов и амплитуд вне массовой поверхности.

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7. Структура основного состояния и несохранение фермионного и барионного чисел в калибровочных теориях

В цикле работ А.Тавхелидзе, выполненных совместно с В.Матвеевым, В.Рубаковым, В.Токаревым, М.Шапошниковым, впервые в рамках стандартной теории электрослабых взаимодействий поставлена и решена проблема нестабильности нормальной барионной материи в экстремальных условиях сверхвысоких плотностей. В этих работах было показано, что ключевым свойством калибровочных взаимодействий, лежащим в основе вывода о нестабильности сверхплотной барионной материи, является сложная структура вакуума и связанное с ней сильное несохранение фермионных квантовых чисел.

Принципиально важным результатом этих исследований является вывод о возможности существования в природе процессов интенсивного распада нормального вещества в контакте с каплей сверхплотной фермионной материи с мощным выделением энергии. В последнее время возможности подобных явлений активно обсуждаются в связи с запуском нового поколения коллайдеров со сталкивающимися релятивистскими ядрами, а также в связи с поиском «темной» материи во Вселенной.

Одним из характерных проявлений сложной структуры вакуума в теории калибровочных полей является продемонстрированное в этом цикле работ на примере точно решаемой модели Швингера — 2-мерной квантовой электродинамики — вырождение вакуума по фермионному числу и киральности при учете топологически нетривиальных конфигураций калибровочных полей (так называемый двойной θ-вакуум).

На основе точного операторного решения уравнений этой модели найдено физическое пространство и построено основное состояние системы, которое не имеет определенных фермионного и кирального чисел и характеризуется двумя произвольными угловыми параметрами. Условие конфайнмента в данной модели сводится к требованию обращения в нуль в пространстве физических состояний оператора калибровочного (электрического) заряда.

Возбуждения системы описываются скалярным нейтральным полем Σ, для которого установлена связь с градиентно инвариантными локальными токами, построенным из фермионных и калибровочных полей. Решена проблема $U(1)$-аномалии. Изучен как механизм несохранения аксиального тока, так и его следствия, в частности, появление кваркового конденсата $\langle\psi\psi\rangle = 0$ и ненулевой массы скалярной Σ-частицы. Показано, что Σ-частица допускает альтернативное представление либо как кварк-антикварковое состояние, либо как состояние псевдоскалярного глюония.

В работах, выполненных совместно с В.Токаревым, метод цепочек уравнений Н.Боголюбова использован для исследований свойств функций Грина калибровочно-инвариантных составных операторов в модели Швингера, и на этой основе установлена структура основного состояния модели.

А.Тавхелидзе (совместно с Н.Красниковым и В.Кузьминым, 1977 г.) в рамках теорий «великого объединения» предложена модель калибровочного взаимодействия со сверхслабым CP-нарушением, позволяющая описать одновременно как эффект CP-нарушения в редких и K-распадах, так и возникновение барионной асимметрии Вселенной.

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