ABOUT STRONG INTERACTION OF FUNDAMENTAL PARTICLES

S.S.SANNIKOV-PROSKURYAKOV

National Scientific Center 'Kharkiv Institute of Physics and Technology“
(*1, Academichna Str., Kharkiv 61108, Ukraine*)

UDC 539.12
‰ 2002

Here, by proceeding from the previous consideration of the problem of interaction [1], we concentrate upon the main properties of strong interaction of hadrons. It is demonstrated that, due to the unusual character of the field propagator in a fiber (at very small distances) where strong interaction is switched on, a new symmetric Green function is used as a field propagator. As a result, the unitary scattering matrix of strong interaction is represented as a T_s-time ordered chronological exponent. It is shown that the particle skeleton algebra plays an important role in finding the full interaction Lagrangian. Coupling constants of strong interactions are determined. They are fixed by the *a priori* standardization of fundamental particle fields. In connection with this, the principle of associative particle creation is formulated. The puzzle of charmed particle existence is guessed. In Appendix, the radiative corrections to the nucleon mass and the masses of η, π, K mesons transfering the strong interactions are calculated.

1. Introduction

Strong interaction of hadrons is apparently the least grounded and comprehended kind of particle interactions. Unification connected with the quark model turns out non-unique and non-obligatory. In the approach to the elementary particle problem suggested in [2], strong interaction has non quark-gluon nature. Strictly speaking, interactions in this theory are appeared only at the particle level; at the level of particle constituents, there are no interactions at all (with the exception of the gravitational one).

We recall that the existence of fundamental particles in this theory is connected with a new dynamical structure ‒ bi-Hamiltonian one by which a new physical reality ‒ ether ‒ is described [2]. Hereby, particles originate in the quantum transition $f \to \dot{f}$ taking place in the system. Degeneration of the ground state of this system is the nearest reason of the existence of the so-called gauge interactions between fundamental particles, see [1]. They are: 1st kind strong (see Section 2), electromagnetic [3], and gravitational [4]. (As regards to the weak interaction, it has another non-gauge nature and is caused by correlation between fermion-fermion fibers [5].) Besides the 1st kind, there is also the 2nd kind strong interaction which is also caused (like weak one) by fiber correlation. But, in this case, they are between fermion-boson fibers (see Section 5).

The basic distinction of strong interaction from electromagnetic and gravitational ones roots in the following: strong interaction is switching on at the very small distances (in the fiber, it is as if capsulated), where there is no space-time yet and therefore a pure field evolution is impossible. There, another kind of motion ‒ involution ‒ plays an important role. As for the electromagnetic interaction, it is switched on in the space-time [3]. The following facts are connected with this: a) the local parameters $\theta(x), \vec{\theta}(x)$ of the strong degeneration group $U_i(1) \otimes \mathrm{SU}_i(2)$ are pseudoscalar immediately observable fields behind which η, π mesons stand (Section 2), while the local parameter $\chi(X)$ of the electromagnetic degeneration group $U_l(1)$ is scalar immediately non-observed field, b) 1st kind strong interactions depend on the helicity of fundamental spinor fields, while the electromagnetic one does not.

As a consequence of this distinction, a new kind of chronological ordering, namely symmetric T_s, appears (Section 3)[1]. Only T_s-operation makes the S-matrix of strong interactions (transferred by non-zero mass particles and used the strong form-factor) to be unitary operator (see Appendix, where the masses of η, π, K mesons as well as the correction to the nucleon mass are calculated). The usual Wick's T_W-operations applicable only in the case of electromagnetic and gravitational interactions transferred by zero mass particles and described by the electromagnetic formfactor (see [7]).

[1] Note that the symmetric Green function (Section 3) connected with this operation as a half sum of retarded and advanced Green functions was actually used by Dirac in the classical theory of electron [6]

In the frame of the suggested scheme, there is an interesting explanation of the origin of so-called charmed particles and charm quantum number (Section 6).

2. Switching on Strong Interaction. Equations of Interacting Fields

a) In the suggested theory, the existence of strong interaction (like electromagnetic and gravitational ones) is connected with degeneration of the ground state $\dot{f}_z$ of the relativistic bi-Hamiltonian system based on the Heisenberg algebra $h^{(*)}_{16}$. This property is described by the degeneration group $U_i(1) \otimes SU_i(2)$ (see its definition in [1]). As (by definition) strong interaction is switched on at supersmall distances (in the hadron era, i.e., in the fiber where there is no space-time almost and therefore the unlimited field evolution is impossible), another kind of motion, namely involution (complex conjugation), plays an exclusively important role. This operation allows one to get (at the skeleton level) the R-spinor χ' from the L-spinor φ by means of the formula $\chi' = -\dot{\varepsilon}\overline{\varphi}$ (recall that the field evolution allows one to obtain the R-field χ from the L-field φ by means of the equality $\chi = -\frac{1}{m} i\sigma^+_\mu \frac{\partial}{\partial X_\mu}\varphi)$.

If the L-spinor φ has fermionic charge $F = 1$ and hypercharge $Y = 1$, and it is transformed under the representation $\left(\frac{1}{2}, 0\right) \otimes D\left(\frac{1}{2}\right)$ of the group $SL_l(2, \mathbf{C}) \otimes SU_i(2)$, so the R-spinor χ' is transformed under the complex conjugate representation $\left(0, \frac{1}{2}\right) \otimes \overline{D}\left(\frac{1}{2}\right)$ of this group and has $F = -1$ and $Y = -1$. Therefore, if $\varphi_{\alpha k} \to \upsilon^{\beta}_{\alpha}\varphi_{\beta m} u^T_{mk}$ ($\upsilon \in SL_l(2, \mathbf{C})$, $u \in SU_i(2)$), where $\varphi_{\alpha k} = \begin{pmatrix} \varphi_{11} & \varphi_{12} \\ \varphi_{21} & \varphi_{22} \end{pmatrix}$ is a square matrix, so $\chi'^{\dot\beta}_{k} \to (u^{-1})^T_{km}\chi'^{\dot\alpha}_{m}\dot\upsilon^{\dot\beta}_{\dot\alpha}$ ($\dot\upsilon = (\upsilon^+)^{-1}$). By introducing the Schwinger's operation to read formulas from the right to the left, i.e., $(u^{-1})^T\chi'\dot\upsilon' \overset{\prime}{\to} \dot\upsilon\chi'(u^{-1})^T$, we see that χ' is transformed under isotopic transformations by means of the inverse matrix u^{-1}.

If spinors $\overline{\varphi}$ could exist in a fiber in the free state like φ ($\overline{\varphi}_{\alpha k}$ are always pairing with $\varphi_{\beta m}$ due to the 100 percent fermion-antifermion asymmetry of the scheme, see [2]), so the skeleton χ' would lead in the quantum transition $f \to \dot{f}$ to the fundamental positive-frequency R-field of antifermion $\chi'(X)$, and together two 2-spinor fields would give a bispinor field $\tilde\psi = \begin{pmatrix} \varphi \\ \chi' \end{pmatrix} = \begin{pmatrix} \varphi \\ \varphi^c \end{pmatrix}$.

The bispinor $\tilde\psi$ has the property: $\tilde\psi^c = \tilde\psi$, where $\tilde\psi^c = C^{-1}\tilde{\overline{\psi}}^T$ is a charge conjugated bispinor and $\tilde{\overline{\psi}} = \tilde\psi^+\gamma_4$ is the Dirac conjugated bispinor. Here, $C = \begin{pmatrix} \varepsilon^{-1} & 0 \\ 0 & -\dot\varepsilon^{-1} \end{pmatrix}$ is the matrix of charge conjugation and $\varepsilon = \dot\varepsilon = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}$. Hence, $\tilde\psi$ is a Majorana bispinor.

At the evolution described by the operator $\gamma_\mu \frac{\partial}{\partial X_\mu}$ (here, $\gamma_\mu = \begin{pmatrix} 0 & -i\sigma^-_\mu \\ i\sigma^+_\mu & 0 \end{pmatrix}$ are the Dirac matrices in the Weyl representation and $\sigma^\pm_\mu = (\vec\sigma, \pm i)$), the field $\tilde\psi$ gives rise to a new field $\tilde\psi'$:

$$\gamma_\mu \frac{\partial}{\partial X_\mu}\tilde\psi = -m\tilde\psi' = -m\begin{pmatrix} \varphi' \\ \chi \end{pmatrix} \quad (1)$$

or $i\sigma^+_\mu \frac{\partial}{\partial X_\mu}\varphi = -m\chi$, $i\sigma^-_\mu \frac{\partial}{\partial X_\mu}\chi' = m\varphi'$ (in our terminology, the field $\tilde\psi'$ as a result of evolution is not fundamental). $\tilde\psi'$ like $\tilde\psi$ is the Majorana bispinor. It is clear that χ (like φ) has $F = Y = 1$ and φ' (like χ') has $F = Y = -1$. As $i\sigma^-_\mu \frac{\partial}{\partial X_\mu}\chi = m\varphi$, $i\sigma^+_\mu \frac{\partial}{\partial X_\mu}\varphi' = -m\chi'$. So, we have

$$\gamma_\mu = \frac{\partial}{\partial X_\mu}\tilde\psi' = -m\tilde\psi. \quad (2)$$

Let us consider the 8-spinor $\tilde\Psi = \begin{pmatrix} \tilde\psi \\ \tilde\psi' \end{pmatrix}$ and 8×8-matrices $\tilde\Gamma_\mu = \begin{pmatrix} \gamma_\mu & 0 \\ 0 & \gamma_\mu \end{pmatrix}$. Then Eqs. (1) and (2) may be written in the form

$$\tilde\Gamma_\mu \frac{\partial}{\partial X_\mu}\tilde\Psi = -mV_0\tilde\Psi,$$

where $V_0 = \begin{pmatrix} \mathbf{0} & \mathbf{1} \\ \mathbf{1} & \mathbf{0} \end{pmatrix}$ and $\mathbf{1}, \mathbf{0}$ are the unit and zero 4×4-matrices (it is obvious that V_0^2 is the unit 8×8-matrix).

Denote $V\tilde\Psi = \Psi$, where

$$V = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \\ 0 & 1 & 0 & 0 \end{pmatrix} \quad (3)$$

and 1, 0 are the unit and zero 2×2-matrices. We have $V^2 = 1$. As $\tilde{\Psi} = \begin{pmatrix} \varphi \\ \chi' \\ \varphi' \\ \chi \end{pmatrix}$ so $V\tilde{\Psi} = \Psi = \begin{pmatrix} \varphi \\ \chi \\ \varphi' \\ \chi' \end{pmatrix} = \begin{pmatrix} \varphi \\ \chi \\ \chi^c \\ \varphi^c \end{pmatrix} = \begin{pmatrix} \psi \\ \psi^c \end{pmatrix}$, where ψ is a usual Dirac bispinor. It is not difficult to verify that $VV_0\tilde{\Gamma}_\mu V^{-1} = \begin{pmatrix} \gamma_\mu & 0 \\ 0 & \gamma_\mu \end{pmatrix} = \Gamma_\mu$. Therefore, we have $\Gamma_\mu \frac{\partial}{\partial X_\mu}\Psi = -m\Psi$, i.e., we come to the usual Dirac equation for the Dirac bispinor field

$$\gamma_\mu \frac{\partial}{\partial X_\mu}\psi = -m\psi \qquad (4)$$

(hereby, $\gamma_\mu \frac{\partial}{\partial X_\mu}\psi^c = -m\psi^c$).

Now we can be convinced that

$$[\Gamma_5, V] = 0,$$

where $\Gamma_5 = \begin{pmatrix} \gamma_5 & 0 \\ 0 & \gamma_5 \end{pmatrix}$ and $\gamma_5 = \gamma_1\gamma_2\gamma_3\gamma_4$ is the Dirac matrix. This means that, under isotopic transformations, $\tilde{\Psi}$ and Ψ are transformed identically: if $\tilde{\psi} \to e^{i\gamma_5(\theta + (1/2)\vec{\tau}\vec{\theta})}\tilde{\psi}$, so

$$\psi \to e^{i\gamma_5(\theta + (1/2)\vec{\tau}\vec{\theta})}\psi. \qquad (5)$$

Thus, if $\varphi \to e^{i(\theta + (1/2)\vec{\tau}\vec{\theta})}\varphi$, so $\chi \to e^{-i(\theta + (1/2)\vec{\tau}\vec{\theta})}\chi$, i.e., χ is as if 'infected" by χ' and χ is transformed by means of the inverse isotopic matrix u^{-1} like χ'. All this results in the statement: *strong interaction depends on the helicity of a spinor field* (matrix γ_5). This fact is exclusively connected with the circumstance that this kind of interaction is switched on inside the fiber (at very small distances), where there is no space-time yet and involution plays a decisive role.

It is seen that, in the Majorana fields $\tilde{\psi}$ and $\tilde{\psi}'$, barions and antibarions are mixed up while they are separated in the Dirac fields ψ and ψ^c. Now fermion-antifermion asymmetry forces to throw away the ψ^c as not arisen in the quantum transition $f \to \dot{f}$. But these fields will arise after switching on interactions.

The general scheme of switching on interaction is considered in [1]. From there and (5), it follows that the degeneration parameters θ and $\vec{\theta}$ must be pseudoscalar magnitudes. They themselves are observable entities (pseudovector current is not conserved) connected with the pseudoscalar fields of η and $\vec{\pi}$ mesons transferring this interaction (unlike the electromagnetic and gravitational ones characterized by conserved currents, see [3, 4]).

So, it is not needed to climb down at the hypothetical quark-gluon level and accept the hypothesis about a composite structure of η and π mesons in order to explain the pseudoscalar character of these particles.

Concerning quarks, it could be said that they might be connected with the model of the bi-Hamiltonian system based on the $h_{24}^{(*)}$ Heisenberg algebra in which isotopic symmetry is represented by the SU(3) group. However, quarks in this model would be observable objects like usual fundamental particles. As quarks are not indeed observable entities, we consider the model $h_{24}^{(*)}$ to be non-realistic and the quark hypothesis to be not needed[2]. From a pure mathematical point of view, the algebra $h_{16}^{(*)}$ is preferred in comparison with $h_{24}^{(*)}$ because the first is connected with dispinors $\varphi_{\alpha k}$ ($\alpha, k = 1, 2$) represented by square matrices and with the algebra of matrices $\sigma_\mu \otimes \tau_m$ having no divisors of zeros (in $h_{24}^{(*)}$, the variables $\varphi_{\alpha,k}$, $\alpha = 1, 2$, $k = 1, 2, 3$ are represented by rectangular matrices). Hereby, in the suggested scheme the 'colour" quantum number does not appear at all.

b) Modulated fields (now we consider only Dirac fields) $\psi^\Sigma(X; Y, Y') = e^{-i\gamma_5\Phi^\Sigma(X,Y')}\psi^\Sigma(X;Y)$ obey the equation of interacting fields (see [1])

$$\gamma_\mu \frac{\partial}{\partial X_\mu}\psi^\Sigma(X) + m\psi^\Sigma(X) =$$

$$= -i\gamma_\mu\gamma_5 \frac{\partial\overbrace{\Phi^\Sigma(X)}}{\partial X_\mu}\psi^\Sigma(X) = -J^\Sigma \qquad (6)$$

(here, the phase Φ is considered to be small and a slow varying function on space-time, that allows one to apply the perturbation method; upper bracket means the integration over inner spatial variables Y, Y', see [1]). We will call gauge 'currents" of the type $J^\Sigma(X) = i\gamma_\mu\gamma_5 \frac{\partial\overbrace{\Phi^\Sigma(X)}}{\partial X_\mu}\psi^\Sigma(X)$ as strong interactions of the 1st kind. Non-Hermitian Lagrangians

$$L_i^\Sigma(X) = \bar{\psi}^\Sigma(X)J^\Sigma(X) \qquad (7)$$

are connected with such 'currents". Such couples arise already inside the fiber and exist between the skeletons

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[2] In connection with cosmology, it is important to notice that another hypothesis about the existence of primary (or relict) black holes (with masses of $10^{-5} \div 10^{13}$ *g*) is not confirmed in the developed scheme: in quantum transitions $f \to \dot{f}$, such objects are not arisen.

of fundamental particles and intermediate mesons (see further).

It is essential that we came to the space-time description of strong interactions in the terms of fundamental particle fields themselves. Therefore with confidence, we can further apply the quantization procedure (in the form of Dirac – Fock second quantization) to the bilocal fields and use Feynman's graphs (generally speaking, it is doubtful to apply this scheme at a deeper level – to the particle constituents, see [2]). We recall that, in the suggested particle theory based on the bi-Hamiltonian dynamical system, hadron fields are bilocal (two-point) functions, see [2] (in the traditional approach based on the quark model, hadron fields are three-point functions).

3. *S*-matrix in Strong Interaction Theory

Proceeding from the non-Hermitian Lagrangian $L_i(X)$ (7) or integrating the equations of interacting fields (6) (comp. with [8]), we can construct the Wick chronological exponent $T_W \exp\,(i\int L_i(X)\,d^4X)$ (where T_W is the Wick chronological ordering). It turns out that this operator will not be unitary even if to work with the Hermitian Lagrangian $\frac{1}{2}(\overline{\psi}J + \overline{J}\psi)$ (and it was not noticed in [1]). The matter is in the form-factor which is used at supersmall distances (when both particles are massive, see [1]). Specific calculations (see Appendix) prompt that, in order to get the unitary scattering matrix S in this situation, we have to accept another axiomatics basically distinct from Feynman's one.

As already noticed, fields in the region of supersmall distances take part in two kinds of motion: continuous evolution and discrete involution. When there is also a space-time, a little particle slightly advancing ahead in time (by means of the operation T_W) is forced to roll back in time (by means of the operation T^+_W; here '+" is involution): strong interaction is as if capsulated in a small region. (It may be said that, in the fiber, the mixture of evolution and involution leads to the local superposition of direct and inverse in time propagation.) As a result, the symmetric operation $T_s = \frac{1}{2}(T_W + T^+_W)$ appears. Therefore, the S-matrix in the theory of strong interaction is written as

$$S = T_s \exp\,(i\int L_i(X)\,d^4X). \tag{8}$$

This operator is unitary (see Appentix), and L_i is automatically an Hermitian Lagrangian.

At disentanglement of S-matrix (8), symmetric pairings appear instead of usual chronological pairing, Feynman's causal Green functions (comp. with [1]):

$$T_s(\psi^\Sigma(X)\,\overline{\psi}^\Sigma(X')) = d^\Sigma\left(\frac{\partial}{\partial X}\right)\Delta^s(X - X'), \tag{9}$$

where Δ^s is the symmetric Green function of a local scalar field determined by formula (see [9], where there is a mistake; P stands for principal value):

$$\Delta^s(X) = \frac{P}{(2\pi)^4}\int \frac{e^{ipX}}{p^2 - M^2}\,d^4X =$$

$$= \frac{1}{4\pi}\left(\delta(X^2) - \theta(X^2)\frac{M}{2\sqrt{X^2}}J_1(M\sqrt{X^2})\right) \tag{10}$$

(this function is zero outside the light cone). For bilocal fields, we obviously have

$$T_s(\psi^\Sigma(X;Y)\,\overline{\psi}^\Sigma(X';Y')) =$$

$$= d^\Sigma\left(\frac{\partial}{\partial X}\right)\Delta^s(X - X';Y,Y'), \tag{11}$$

where

$$\Delta^s(X;Y,Y') = F\left(Y; i\frac{\partial}{\partial X}\right)F\left(Y';\frac{\partial}{\partial X}\right)\Delta^s(X) \tag{12}$$

(here, $F(Y; i\frac{\partial}{\partial X})$ is a smearing operator, see [2]). By definition, we have

$$\Delta^s = \frac{1}{2}(\Delta^c + \Delta^a) = \frac{1}{2}\varepsilon(X_0)\,\Delta = \frac{1}{2}(\Delta^{\rm ret} + \Delta^{\rm adv})$$

(the latter form was used by Dirac [6]), where Δ^c, Δ^a are causal and anticausal Green functions (Δ^a is complex conjugate to the Δ^c, $\varepsilon(X_0)$ – the sign-function, Δ – permutation Pauli – Jordan function). The following relations take place for Δ^c:

$$\Delta^c = \frac{1}{2}(\varepsilon(X_0)\,\Delta + i\,\Delta'),\quad \Delta^c - \Delta^a = i\,\Delta'$$

(Δ' is the antipermutation function). In connection with this, it is interesting to consider a new Green function $\tilde{\Delta}$ and its complex conjugate $(\tilde{\Delta})^+$: $\tilde{\Delta} = \frac{1}{2}(\Delta + i\,\varepsilon(X_0)\,\Delta') = \varepsilon(X_0)\,\Delta^c$ hereby $\tilde{\Delta} + (\tilde{\Delta})^+ = \Delta$ and

$\tilde{\Delta} - (\tilde{\Delta})^{+} = i\,\varepsilon\,(X_0)\,\Delta'$. Then we have else $\Delta^{s} = \frac{1}{2}\,\varepsilon\,(X_0)\,(\tilde{\Delta} + (\tilde{\Delta})^{+})$.

It is clear that, in the frame of this new axiomatics, we completely loss the possibility to use the dispersion relations method [9] for description of strong interaction processes (we note that 'currents“ $J\,(X)$ are local only in form; in fact, they are essentially non-local).

4. Identities in the Skeleton Algebra. Currents, Lagrangians. Coupling Constants of Strong Interactions

Let L be a Lagrangian plane with coordinates $\varphi_{\alpha k}, \bar{\varphi}_{\dot{\alpha} k}$ $(\alpha, k = 1, 2)$ in the algebra $h^{(*)}_{16}$. In the suggested theory, the non-Fock representation of $h^{(*)}_{16}$ is considered. Such a representation contains additional variables $-$ Lorentzian scalar $\varphi_k \doteq \varphi_{2k}$, $\bar{\varphi}_k \doteq \bar{\varphi}_{2k} \in L_0$ (they are the proper variables of ether); therefore, $L_0 \subset L$. Let us consider the commutative graduated algebra $U\,[L]$ over L. It is the maximal commutative subalgebra in the envelope algebra $U\,[h^{(*)}_{16}]$. $U\,[L]$ is our skeleton algebra.

a) I d e n t i t i e s. In the ground of the identities mentioned above between the skeletons of fundamental particles O^{Σ} as elements of the algebra $U\,[L]$ and parameters $\eta, \vec{\pi}$ of the degeneration group $U_i\,(1) \otimes \mathrm{SU}_i\,(2)$ playing an important role in the origination of possible 1st kind couples (Lagrangians, see further), the completeness condition between the 2×2-matrices $\tau_m = (\vec{\tau}, 1)$ lines:

$$1_{km}\,1_{np} + \vec{\tau}_{km}\,\vec{\tau}_{np} = 2\delta_{kp}\,\delta_{nm}. \qquad (13)$$

Pairing up this condition with the additional variables $\bar{\varphi}_k\,\varphi_p$, we obtain the formula (condition for out-pairing)

$$2\,\bar{\varphi}\,\varphi\,\delta_{nm} = \varphi_n\,\bar{\varphi}_m + (\vec{\tau}\varphi)_n\,(\bar{\varphi}\,\vec{\tau})_m, \qquad (14)$$

in which the pairing variables $\bar{\varphi}\,\varphi = \bar{\varphi}_k\,\varphi_k$ are expressed by free ones $\varphi_n\,\bar{\varphi}_m$. If to wind up formula (13) with $\varphi_{\alpha m}\,\eta$, where $\varphi_{\alpha m} = N_{\alpha m}$ is the skeleton of a nucleon (α are the Dirac's indices) and η is a parameter of the group $U_i\,(1)$, the skeleton of η-meson, we get the first identity for $N = 3$

$$2\bar{\varphi}\,\varphi\,\varphi_{\alpha n}\,\eta = 2N^{*}_{\alpha n}\,\eta = \bar{\varphi}_m\,\varphi_{\alpha m}\,\varphi_n\,\eta +$$

$$+\ \bar{\varphi}\,\vec{\tau}\varphi_{\alpha}\,(\vec{\tau}\varphi)_n\,\eta = a\,(\Lambda_{\alpha}\,K_n + \vec{\Sigma}_{\alpha}\,(\vec{\tau}K)_n\,) \qquad (15)$$

(further we limit ourselves by identities linear in η and $\vec{\pi}$ only). In (15), $N^{*}_{\alpha n} = \bar{\varphi}\,\varphi\,\varphi_{\alpha n}$, $\Lambda_{\alpha} = \bar{\varphi}\,\varphi_{\alpha}$, $\vec{\Sigma}_{\alpha} = \bar{\varphi}\,\vec{\tau}\varphi_{\alpha}$ are the skeletons of the first nucleon resonance, Λ and Σ $-$ hyperons, and $\eta\,\varphi_n = aK_n$ is the skeleton of new K mesons which are created when the additional variables φ_n join the η $-$ skeleton. (As a result, a pseudoscalar isotopic dublet arises from a pseudoscalar isotopic singlet; a is a real number which we can put at once by definition to be equal to 1).

Further, let condition (14) be wound up with $(\hat{\pi}\,N_{\alpha})_n$, where $\hat{\pi} = \vec{\pi}\vec{\tau}$ ($\vec{\pi}$ are the skeleton of π mesons). Using the well-known relation

$$\vec{\tau}\hat{\pi} = \vec{\pi} + i\,[\vec{\pi}, \vec{\tau}] \qquad (16)$$

and the formula following from it,

$$\vec{\pi}\varphi = b\,\vec{\tau}K + i\,[\vec{\tau}, \vec{\pi}]\,\varphi, \qquad (17)$$

where we put else $\hat{\pi}\,\varphi = bK$ (b is a real number), we get

$$2\,(\hat{\pi}\,N^{*}_{\alpha})_n = b\,(\Lambda_{\alpha}\,K_n + \vec{\Sigma}_{\alpha}\,(\vec{\tau}K)_n\,) + 2\,i\,[\vec{\pi}, \vec{\Sigma}_{\alpha}]\,(\vec{\tau}\varphi)_n. \qquad (18)$$

Since by definition $K = \frac{1}{a}\,\eta\,\varphi = \frac{1}{b}\,\hat{\pi}\,\varphi$, this implies

$$(\hat{\pi} - \frac{b}{a}\,\eta)\,\varphi = 0. \qquad (19)$$

Multiplying the last equality by $\hat{\pi} + \frac{b}{a}\,\eta$, we obtain

$$(\vec{\pi}^{2} - \varepsilon^{2}\,\eta^{2})\,\varphi = 0,$$

where $\varepsilon = \frac{b}{a}$ ($\varepsilon = b$, if $a = 1$).

Futher, it follows from (16) that

$$\varepsilon\,\eta\,\bar{\varphi}\,\vec{\tau}\varphi = \vec{\pi}\,\bar{\varphi}\,\varphi + i\,[\vec{\pi}, \bar{\varphi}\,\vec{\tau}\varphi],$$

i.e., the following two equalities are true:

$$\vec{\pi}\,\bar{\varphi}\,\varphi = \varepsilon\,\eta\,\bar{\varphi}\,\vec{\tau}\varphi, \quad [\vec{\pi}, \bar{\varphi}\,\vec{\tau}\varphi] = 0.$$

Therefore, it has to consider $\vec{\pi}$ and $\bar{\varphi}\,\vec{\tau}\varphi$ to be co-linear isovectors. However, $\vec{\pi}$ and $\bar{\varphi}\,\vec{\tau}\varphi$ are magnitudes from different vector spaces: $\vec{\pi}$ are parameters of the degeneration group, but $\bar{\varphi}\,\vec{\tau}\varphi$ are dynamical variables of the system. Hence, we may not put $\vec{\pi} = \kappa\,\bar{\varphi}\,\vec{\tau}\varphi$ and by analogy $\eta = \kappa'\,\bar{\varphi}\,\varphi$, where κ and κ' are any real numbers connected by the condition $\kappa = \varepsilon\,\kappa'$. Otherwise, we would have

$$\vec{\pi}^{2} - \eta^{2} = \kappa^{2}(\bar{\varphi}\,\vec{\tau}\varphi)^{2} - \kappa'^{2}(\bar{\varphi}\,\varphi)^{2} = \kappa^{2}(1 - \varepsilon^{-2})(\bar{\varphi}\,\varphi)^{2},$$

as $(\bar\varphi\ \vec\tau\varphi)^2 = (\bar\varphi\varphi)^2$ $((\bar\varphi\ \vec\tau\varphi, i\bar\varphi\varphi)$ is an isotropic 4-vector). But it is impossible, as $\vec\pi^2 - \eta^2$ is not a dynamical variable ($(\bar\varphi\varphi)^2$ is such a variable). A solution of this trouble is seen in putting $\varepsilon^2 = 1$, i.e., $b = \pm a$ (and if $a = 1, \varepsilon = b = \pm 1$). Then the bond between $\vec\pi$ and η has the form $(\vec\pi^2 - \eta^2)\varphi = 0$ or $\vec\pi^2 - \eta^2 = 0$. This means that the standardization of π mesons is equal to the η meson one.

Number b may be connected with the G-parity of η and π mesons. As G-parity of η and π are opposite, we have to choose $b = -1$ (however further, we will consider b and ε to be indefinite magnitudes).

It follows from (15), (18) that

$$((\hat\pi - \varepsilon\eta) N^*_\alpha)_n = i[\vec\pi, \vec\Sigma_\alpha](\vec\tau\varphi)_n. \tag{20}$$

Generally speaking, $(\hat\pi - \varepsilon\eta)\varphi_\alpha \neq 0$ although $(\hat\pi - \varepsilon\eta)\varphi = 0$, where $\varphi_n \doteq \varphi_{2n}$. It means that there are no identities at $N = 1$.

Pairing (20) with $\bar\varphi_n$ and taking into account that $[\vec\pi, \vec\Sigma_\alpha]\ \bar\varphi\ \vec\tau\varphi = [\bar\varphi\ \vec\tau\varphi, \vec\pi]\ \vec\Sigma_\alpha = 0$, we obtain $\bar\varphi\hat\pi N^*_\alpha = \varepsilon\eta\bar\varphi N^*_\alpha$. As $\bar\varphi\hat\pi N^*_\alpha = \bar\varphi\varphi\vec\Sigma_\alpha\vec\pi$ and $\bar\varphi N^*_\alpha = \bar\varphi\varphi\Lambda_\alpha$, we have ($N = 2$)

$$\vec\Sigma_\alpha\vec\pi = \varepsilon\eta\Lambda_\alpha. \tag{21}$$

Multiplying (17) on $\vec\Sigma_\alpha$ and using (21), we come to the identity

$$i[\vec\pi, \vec\Sigma_\alpha](\vec\tau\varphi)_n = b(\Lambda_\alpha K_n - \vec\Sigma_\alpha(\vec\tau K)_n). \tag{22}$$

Hence, (18) may be written in the form ($N = 3$)

$$2\hat\pi N^*_\alpha = b(3\Lambda_\alpha K - \hat\Sigma_\alpha K),\quad \hat\Sigma_\alpha = \vec\Sigma_\alpha\vec\tau, \tag{23}$$

and, in particular, we have $(\hat\pi - \varepsilon\eta)N^*_\alpha = b(\Lambda_\alpha - \hat\Sigma_\alpha)K$. By multiplying (20) from the left by $\hat\pi + \varepsilon\eta$, we obtain $(\vec\pi^2 - \eta^2)N^*_\alpha = 0$ as $(\hat\pi - \varepsilon\eta)\varphi = 0$ and $[\vec\pi, \vec\Sigma_\alpha]\ \vec\pi = 0$.

Winding up the complex conjugate relation to (17), $\bar\varphi_n\vec\pi = b(\bar K\vec\tau)_n - i(\bar\varphi[\vec\tau, \vec\pi])_n$, with $\varphi_{\alpha n}$, we obtain: $\Lambda_\alpha\vec\pi = b\bar K\vec\tau N_\alpha - i[\vec\Sigma_\alpha, \vec\pi]$. At $N = 2$, we have

$$[\vec\Sigma_\alpha, \vec\pi] = i(\Lambda_\alpha\vec\pi - b\bar K\vec\tau N_\alpha). \tag{24}$$

Obviously,

$$\bar K\vec\tau N_\alpha = \frac{1}{a}\eta\vec\Sigma_\alpha. \tag{25}$$

From here, by multiplying by $\vec\pi$, we get (at $N = 2$)

$$\vec\Sigma_\alpha\vec\pi = b\bar K N_\alpha. \tag{26}$$

From (21) and (26), it follows that (at $N = 2$)

$$\varepsilon\eta\Lambda_\alpha = b\bar K N_\alpha. \tag{27}$$

Therefore at the same time, we have (at $N = 2$)

$$[\vec\Sigma_\alpha, \vec\pi] = i(\Lambda_\alpha\vec\pi - \varepsilon\vec\Sigma_\alpha\eta). \tag{28}$$

In a quite analogous manner, we can obtain the identities for $N = 4$ and $N = 5$ between $\Xi_{\alpha n}$, Λ^*_α, $\vec\Sigma^*_\alpha$ hyperons and $\eta, \vec\pi, K$ mesons. However, we will not stop here at these identities.

b) C u r r e n t s a n d L a g r a n g i a n s. Identities (15), (21), (23) – (29) play an important role in the finding of couples between the fields of fundamental particles and η, π, K mesons – Lagrangians of the 1st kind strong interactions. At switching on these interactions, the equations $(\Gamma^\Sigma_\mu \frac{\partial}{\partial X_\mu} + M_\Sigma)\Psi^\Sigma = 0$ for free fields are transformed into the equations for interacting fields $(\Gamma^\Sigma_\mu \frac{\partial}{\partial X_\mu} + M_\Sigma)\Psi^\Sigma = -J^\Sigma$, see [1]. In the case of N and Λ, Σ hyperons (we consider only these low lying levels in the baryon spectrum), we have for 'persecute" currents:

$$J^N = i\gamma_\mu\gamma_5\left(\frac{\partial}{\partial X_\mu}\eta + \frac{1}{4}\frac{\partial\hat\pi}{\partial X_\mu}\right)N = i(\eta + \frac{1}{4}\hat\pi)N,$$

$$J^\Lambda = 0,$$

$$\vec J^\Sigma = \frac{1}{2}\gamma_\mu\gamma_5\left[\vec\Sigma, \frac{\partial\vec\pi}{\partial X_\mu}\right] = \frac{1}{2}[\vec\Sigma, \vec\pi] \tag{29}$$

(further, spinor indices and space-time derivatives are omitted).

Now we consider couples between particle skeletons only (at the skeleton level). As a result of the transition $f \to \dot f$, the skeletons O^Σ grow over by dependence on space-time coordinates X_μ and Y_μ, i.e., turn out into fields $\psi^\Sigma(X; Y)$ with the *a priori* standardization N_Σ, see [2] (for brevity, we denote $O^\Sigma_i N_\Sigma$ as N_Σ), and couples between skeletons will give the couples between particle fields.

Adding the corresponding identities given above (i.e. 'zeros") to couples (29), we may write

$$J^N = i(\eta + \frac{1}{4}\hat\pi)N + iA[\varepsilon\eta N^* - \frac{b}{2}(\Lambda K + \hat\Sigma K)] +$$

$$+\ i B[\hat{\pi} N^* - \frac{b}{2}(3\Lambda K - \hat{\Sigma} K)],$$

$$J^{\Lambda} = i C(b\bar{K} N - \varepsilon\Lambda\eta) + i D(\vec{\pi}\vec{\Sigma} - \varepsilon\eta\Lambda),$$

$$\vec{J}^{\Sigma} = \frac{1}{2}[\vec{\Sigma}, \vec{\pi}] + E([\vec{\Sigma}, \vec{\pi}] - i\vec{\pi}\Lambda + i\varepsilon\eta\vec{\Sigma}) +$$

$$+\ i F(b\bar{K}\vec{\tau}N - \varepsilon\eta\vec{\Sigma}). \quad (30)$$

Here, A, B, C, D, E, F are any real numbers.

At the skeleton level, these expressions will lead to the Hermitian couples in whole if these numbers obey the conditions

$$\varepsilon A = 1,\ B = \frac{1}{4},\ \frac{1}{2}(A + 3B) = C,$$

$$\frac{1}{2}(A - B) = F,\ D = E. \quad (31)$$

As $\varepsilon = -1$, we have $C = -\frac{1}{8}$, $F = -\frac{5}{8}$.

Although, at first sight, nothing changes at the skeleton level (zeros were added), but after the transition $f \to f$ gone in the relativistic bi-Hamiltonian system, new couples, currents, and Lagrangians of particle fields will arise. From the SU(3) ‾ group point of view, they are F-type couples, see [10]. Note that the couples $\bar{N}\eta N^*$ and $\bar{N}\hat{\pi}N^*$ indeed exist, see [11]. Obviously, due to the usage of involution operation, + (see Section 3), the Lagrangian is hermitized.

Particle skeletons O^{Σ} are graduated by *isoton* number N (for example, $\eta, \vec{\pi}$ are skeletons with $N = 0, N_{\alpha}$; K^- with $N = 1$; $\Lambda, \vec{\Sigma}^-$ with $N = 2$ and so on). At switching on the strong interaction, more and more high values of N will be taken into account. It is obvious that couples of the 1st kind are not closed. In order to get results in the frame of the S-matrix formalism, it is necessary to cut an infinite chain of couples. It is natural to connect this cutting with the fact of smallness of the probability of large N (N obey the Poisson distribution, see [2]) and also of smallness of the *a priori* standardization N_{Σ} for large N. Thus, we come to the so-called Tamm ‾ Dankov method (comp. with the Heisenberg theory [12]). But in our theory (unlike the [12]), this method leads due to the just now mentioned reason to the well-convergent results.

c) Coupling constants. We have to keep in mind that the particle fields in (30) contain the *a priori* standardization N_{Σ} as well as the fields of η, π, K mesons. They satisfy the conditions $N_{N^*} < N_{\Sigma} < N_{\Lambda} << N_N$. Hereby, the nucleon standardization $|N_N| = |N_p| \approx |N_n|$ is equal $\frac{1}{\sqrt{2}}$ as $|N_p|^2 + |N_n|^2 = 1$, see [2].

Now we would like to determine the *a priori* constant f (standardization of η and π mesons). In zero approximation, η and π are connected with nucleons only: N_N is largest. We may close the field system $\{N, \eta, \pi\}$ adding the Lagrangian $L_M = -\frac{1}{4\pi}\frac{1}{2}(\eta^2 + \vec{\pi}^2)$ of free mesons η and π to the Lagrangian $L_N + L_i = \bar{N}N + if\bar{N}(\eta + \frac{1}{4}\hat{\pi})N$. We write the nucleon Lagrangian on the whole and in the exact form as

$$L(-\infty) = |N_N|^2\left[\bar{N}DN + if\bar{N}\gamma_{\mu}\gamma_5 \times \right.$$

$$\left.\times\left(\frac{\partial}{\partial X_{\mu}}\widehat{\eta N} + \frac{1}{4}\frac{\partial}{\partial X_{\mu}}\widehat{\hat{\pi} N}\right)\right] + if\frac{N_{N^*}}{N_N}\bar{N}\gamma_{\mu}\gamma_5 \times$$

$$\times\left(\frac{\partial}{\partial X_{\mu}}\widehat{\eta N^*} + \frac{1}{4}\frac{\partial}{\partial X_{\mu}}\widehat{\hat{\pi} N^*}\right) + if\frac{N_{\Lambda}}{8N_N}\times$$

$$\times\ b\bar{N}\gamma_{\mu}\gamma_5\widehat{\Lambda\frac{\partial}{\partial X_{\mu}}K} + if\frac{5N_{\Sigma}}{8N_N}b\bar{N}\gamma_{\mu}\gamma_5\widehat{\hat{\Sigma}\frac{\partial}{\partial X_{\mu}}K} -$$

$$-\frac{f^2}{4\pi}\frac{1}{2}\left[\left(\frac{\partial\eta}{\partial X_{\mu}}\right)^2 + \left(\frac{\partial\vec{\pi}}{\partial X_{\mu}}\right)^2 + 2\frac{\partial\bar{K}}{\partial X_{\mu}}\frac{\partial K}{\partial X_{\mu}}\right], \quad (32)$$

where $D = \gamma_{\mu}\frac{\partial}{\partial X_{\mu}} + M_N$ and $N, ..., \eta, ...$ are the fields normalized in the usual way (factor $\frac{1}{4\pi}$ in L_M-Lagrangian says that we deal with the Gauss unit system, comp. with [7], where the electromagnetic interaction is considered in this system). In order to get the usual Lagrangian $L(in)$ used in the S-matrix formalism (see [1]), we have to put

$$f^2 = |N_N|^2 = \frac{1}{2} \quad (33)$$

(common phase is fixed by the condition $f = N_N$). Then

$$L(in) = \bar{N}DN + iN_N\bar{N}\gamma_{\mu}\gamma_5\left(\frac{\partial}{\partial X_{\mu}}\widehat{\eta N} + \frac{1}{4}\frac{\partial}{\partial X_{\mu}}\widehat{\hat{\pi} N}\right) +$$

$$+\ i N_{N^*}\overline{N}\gamma_\mu\gamma_5\left(\frac{\partial}{\partial X_\mu}\overbrace{\eta N^*}+\frac{1}{4}\frac{\partial}{\partial X_\mu}\overbrace{\hat{\pi} N^*}\right)+$$

$$+\ \frac{i}{8}b\overline{N}\gamma_\mu\gamma_5\left(N_\Lambda\overbrace{\Lambda\frac{\partial K}{\partial X_\mu}}+5N_\Sigma\overbrace{\hat{\Sigma}\frac{\partial K}{\partial X_\mu}}\right)-$$

$$-\ \frac{1}{8\pi}\left[\left(\frac{\partial\eta}{\partial X_\mu}\right)^2+\left(\frac{\partial\vec{\pi}}{\partial X_\mu}\right)^2+2\frac{\partial\overline{K}}{\partial X_\mu}\frac{\partial K}{\partial X_\mu}\right]. \qquad (34)$$

Hence, $|f|=\frac{1}{\sqrt{2}}$. In (34), $f=N_N$ plays the role of coupling constant of the couple $\overline{N}\overbrace{N\eta}$. Thus, the *a priori* standardization of η, π and K mesons is determined by the condition (33) which we call a *principle of associative creation* of η, π (or K), mesons with nucleons N. According to this principle, the constants of the couples $\overline{N}(\overbrace{\eta N^*}+\frac{1}{4}\overbrace{\hat{\pi} N^*})$, $\overline{N}\overbrace{\Lambda K}$, and $\overline{N}\overbrace{\Sigma K}$, as follows from (34), are equal to N_{N^*}, $-\frac{1}{8}N_\Lambda$ and $-\frac{5}{8}N_\Sigma$, correspondingly, i.e., are negligible. Hence, the main part of the nucleon strong interaction Lagrangian looks as

$$L_i(X)=i\frac{1}{\sqrt{2}\,k}\overline{N}(X)\gamma_\mu\gamma_5\times$$

$$\times\left(\overbrace{\frac{\partial\eta(X)}{\partial X_\mu}N(X)}+\overbrace{\frac{1}{4}\frac{\partial\hat{\pi}(X)}{\partial X_\mu}N(X)}\right). \qquad (35)$$

According to the associative principle, the *a priori* standardization constant f depends on the fundamental particle with which η, π, K are connected. So, if η, π, K mesons are connected with Λ-hyperon, the corresponding Lagrangian has the form (see (30))

$$|N_\Lambda|^2[\overline{\Lambda}\Lambda-if\varepsilon(-\frac{1}{8}+D)\overline{\Lambda}\overbrace{\Lambda\eta}-$$

$$-\ i\frac{1}{8}fb\frac{N_N}{N_\Lambda}\overline{\Lambda}\overbrace{N\overline{K}}+ifD\frac{N_\Sigma}{N_\Lambda}\overline{\Lambda}\overbrace{\vec{\Sigma}\vec{\pi}}-$$

$$-\ \frac{f^2}{4\pi}\frac{1}{2}[\eta^2+\vec{\pi}^2+2\overline{K}K], \qquad (36)$$

and the constant

$$f=N_\Lambda<<N_N. \qquad (37)$$

The coupling constant of the couple $\overline{\Lambda}\overbrace{N\overline{K}}$ is equal to $-\frac{1}{8\sqrt{2}}$, others are negligible.

If η, π, K mesons are connected with Σ-hyperon, the corresponding Lagrangian is

$$|N_\Sigma|^2[\overline{\vec{\Sigma}}\vec{\Sigma}+(\frac{1}{2}+D)f\overline{\vec{\Sigma}}[\vec{\Sigma},\vec{\pi}]+if\varepsilon(\frac{5}{8}+D)\overline{\vec{\Sigma}}\vec{\Sigma}\eta-$$

$$-\ if\frac{N_\Lambda}{N_\Sigma}D\overline{\vec{\Sigma}}\Lambda\vec{\pi}-if\frac{5}{8}b\frac{N_N}{N_\Sigma}\overline{\vec{\Sigma}}\overline{K}\vec{r}N]-$$

$$-\ \frac{f^2}{4\pi}\frac{1}{2}[\eta^2+\vec{\pi}^2+2\overline{K}K], \qquad (38)$$

and

$$f=N_\Sigma<<N_N. \qquad (39)$$

The coupling constant of the couple $\overline{\Sigma}\overbrace{N\overline{K}}$ is equal to $\frac{5}{8\sqrt{2}}$, others are negligible[3].

It is interesting to notice that the coupling constant of a pseudovector couple with derivative is indeed a dimensional magnitude $\frac{1}{\sqrt{2}\,k}\sim 10^{-14}$ cm. Therefore, calculations in this theory in the framework of perturbation theory are connected with the quasi-classical decomposition in the constant k^{-1} (so-called WKBJ-method).

5. The Second Kind Couples (Correlation between Fibers)

In Section 2, it was shown how 1st kind couples arise in the case of fundamental baryons. Quite analogously, 1st kind couples arise in the case of fundamental mesons. So, for example, there are Lagrangians of the form $i\overline{\tilde{K}}^*_{\mu\nu}\overbrace{K^*_\nu\frac{\partial}{\partial X_\mu}\eta}$ or $i\overline{\tilde{K}}^*_{\mu\nu}\overbrace{\vec{\tau}K^*_\nu\frac{\partial}{\partial X_\mu}\vec{\pi}}$ where $\tilde{K}^*_{\mu\nu}=\varepsilon_{\mu\nu\rho\sigma}K^*_{\rho\sigma}$ and $K^*_{\rho\sigma}=\frac{\partial K^*_\rho}{\partial X_\sigma}-\frac{\partial K^*_\sigma}{\partial X_\rho}$ (here, K^*_μ is the field of fundamental vector mesons). However, as vector mesons are non-stable, couples of this kind (linear in η and π) are not essential. Another couples of the first kind (bilinear in η and π) and also couples of the second kind are essential in this case.

The 2nd kind couples originated by the fermion-boson correlation between various fibers are already arisen at the skeleton level. Transitions $f\to\dot{f}$ in these fibers lead to the couples between fundamental particle fields (fermions and bosons). So, for example, the

[3] We note that the coupling constants of other kinds of interaction – electromagnetic and gravitational – as interactions switched on in the space-time are determined by another conditions (see [3, 4]). Concerning the Fermi coupling constant of weak interaction, see [5].

ISSN 0503-1265. Ukr. J. Phys. 2002. V. 47, N 7

correlations $\gamma_\mu \overbrace{\vec{\tau} N \otimes \vec{\rho}_\mu}$, $\gamma_\mu \overbrace{N \otimes \omega_\mu}$ and Lagrangians $\overline{N}\,\gamma_\mu\, \vec{\tau} N\, \rho_\mu$, $\overline{N}\,\gamma_\mu\, N\, \omega_\mu$ or $\dot{\overline{\Delta}}\, N\, \rho$ may arise (here, Δ is a resonance $(\frac{3}{2}, \frac{3}{2}))$. Of course, such correlations may not lead to decays.

It is important to note that these couples are not conditioned by any gauge symmetry of fundamental particle Lagrangians, connected with transformations of the type $\omega_\mu \to \omega_\mu + \frac{\partial\,\theta}{\partial X_\mu}$ (the second kind) or $\psi \to e^{i\theta}\,\psi$ (the first kind): due to the non-zero mass of funfamental mesons, such a symmetry is not admitted [4].

With the existence of 2nd kind couples, such a fact is connected: besides the fundamental component (in the sector T^1), the field of, for example, a resonance Δ may have the component in the sector T^2 which is a bound state in the system (N, ρ). It turns out that a particle skeleton in the sector T^1 may be equal to zero at all, and all its components lie in higher sectors. So, for example, the matter stands in the case of Ω^--hyperon. Its skeleton in the sector T^1 of the form $\left(\tau = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}\right)$:

$$O^{\Omega^-} = \gamma_\nu\, \psi_k\, \overline{\varphi}_k\, \overline{\psi}_n\, \gamma_\mu\, \psi_m\, \overline{\varphi}_m\, (\tau\,(\eta + \frac{1}{4}\hat{\pi}\,))^{np}\, \overline{\psi}_p\, \gamma_\nu\, \psi_s\, \overline{\varphi}_s \tag{40}$$

is forbidden by the selection rules. However, Ω^- to be a composed particle (see the definition in [2]) has the skeleton

$$\gamma_\nu\, \psi_k\, \overline{\varphi}_k\, \overline{\psi}_n\, \gamma_\mu\, \psi_m\, \overline{\varphi}_m\, (\tau\,\hat{\pi}\,)^{np} \otimes \overline{\psi}'_p\, \gamma_\nu\, \psi'_s\, \overline{\varphi}'_s \tag{41}$$

(Singlet is indeed impossible because $\overline{z}\,\tau\,\overline{z} = 0$.) (40) might be called a 'shadow“ of (41) in the sector T^1. The mass of Ω^--hyperon given by (36) in [2] (at $N = 8, i = 1$ and $\mu^2 = 0.067$) is equal to $M_{\Omega^-} = 1.59\,\frac{kh}{c}$ that is not far from the experimental value 1.67 GeV/c^2 [11] ($khc = 1$ GeV).

6. Higher Symmetries

In the fiber, there is the fundamental isotopic symmetry $U_i(1) \otimes SU_i(2)$. At the joining of various fibers and forming of composed particles, the possibility of constructing a higher (non-fundamental) symmetry and their multiplets arises. So, by a simple joining of SU(2)-baryons, the SU(3)-octet $(N, \Lambda, \vec{\Sigma}, \Xi)$ is obtained. As a result, the approximate SU(3)-symmetry without quarks (i.e., SU(3)/$\mathbf{Z}_3$) appears. (We recall that SU(3) is considered in the traditional approach to be an exact symmetry of a hypothetical superstrong interaction that is absent in the framework of the suggested scheme; as for the SU(3)-symmetry connected with the Heisenberg algebra $h_{24}^{(*)}$, see above).

It is very important to emphasize that hypercharge Y is already at the fundamental level: it is connected with the $U_i(1)$-group. At the phenomenological SU(3)-level, a new hypercharge Y' may be introduced. Hereby, eigenvalues of Y' on the particle fields must coincide with values of fundamental hypercharge $Y = -\,\overline{\Phi}\,\Phi - 4$, which is the generator of the $U_i(1)$-group acting on the particle skeleton O^{Σ}.

Quite analogously, the matter stands with other higher flavored symmetries, and quantum numbers connected with them, for example, charm quantum number C. The operator (matrix) C (like Y') may be defined only on particle fields ψ^{Σ} if one preliminarily joins SU(2)-multiplets in flavored-multiples.

However, for particle skeletons O^{Σ}, we may not write the formula $T\,O^{\Sigma}\,T^{-1} = U^{\Sigma\Sigma'}\,O^{\Sigma'}$ ($U^{\Sigma\Sigma'}$ is a matrix representation of the flavored group), where T is an element of the dynamical group $\mathrm{Sp}^{(*)}(8, \mathbf{C})$ of the bi-Hamiltonian system because the latter does not contain the flavored group (comp. with parity P', appearing at the particle level, see [1]).

A question arises: what role might charm C play in the ground of the suggested scheme? As Y' and C always appear as the sum $Y' + C$ (hereby, Y' commutes with C), so it is not difficult to understand that the role of C might take on itself only the operator of fundamental hypercharge Y. In such a case in the space of fields, we must define the hypercharge in a new fashion: now, hypercharge is the sum $Y' + C = \mathbf{Y}$. Eigenvalue of $\mathbf{Y}$ on ψ^{Σ} must coincide of course with the eigenvalue of the fundamental hypercharge Y in ψ^{Σ}. As it must be $Y' = B + S'$ and $\mathbf{Y} = B + \mathbf{S}$, we need to define strangeness in a new fashion: now, $\mathbf{S} = S' + C$. Thus, at the particle classification, the number C is not obligatory (like the muon number [5] or leptonic and baryonic charges: the latters are fermionic charges at the leptonic and barionic branches, correspondingly [2]). The same will be for beauty and top.

Then particles Λ_c^+ $(i = 0)$, $\vec{\Sigma}_c$ $(i = 1)$, F^+, F^{+*} $(i = 0)$ must have $Y = 2$. Hereby, baryons $\Lambda_c^+, \vec{\Sigma}_c$ will

[4] We notice that the correlation between fermion-fermion fibers lying in the base of particle weak interactions, see [5], may not be reduced to the gauge transformation of the mentioned type too.

have $S = 1$ and mesons $F^+, F^{+*} - S = 2$. Mesons D and D^* $(i = 1/2)$ have $Y = S = 1$ and J/ψ $(i = 0)$ has $Y = S = 0$.

As is known, charmed particles have large masses. A question is: why are barions $(F = 1)$ with positive strangeness very heavy (have large N) and are there not light baryons with $S = 1$?

The existence of fundamental baryons (in the sector T^1) with positive strangeness is forbidden by selection rules: the pairing of additional variables φ_k with $\psi_{\alpha k}$ is impossible (φ_k are obligatory paired with $\bar{\psi}_{\dot{\alpha} k}$, and $\bar{\varphi}_k$ with $\psi_{\alpha k}$ only, see [1]). Baryons with $S > 0$ might be composed only. In order to obtain a baryon with $S = 1$, it is needed to pair φ_k from one fiber with $\psi_{\alpha k}$ from another fiber. But there are no free φ_k's in the fiber.

It turns out that, in the sector T^4, there exists an unusual ('heavy“) spinor $\chi_k^{\dot{\alpha}} = \bar{\Phi}_k^{\dot{\alpha}}$ with $N = 31$ and the additional variable $\chi_k \doteq \chi_k^2$ connected with it, which may stick to φ_k. Its existence is connected with a special cluster the Lorentzian scalar formed from $\bar{\varphi}_{\dot{\alpha} k}$, $\bar{\varphi}_{\beta m}$. The latter has isotonic quantum number $N = 32$ and arises as a result of correlation between four fibers. Here is this cluster (aggregate):

$$\mathbf{O} = \varepsilon_{\mu\nu\rho\sigma}\varepsilon_{\mu'\nu'\rho'\sigma'}\varepsilon_{\mu''\nu''\rho''\sigma''}\varepsilon_{\mu'''\nu'''\rho'''\sigma'''}\varepsilon_{mm'm''m'''}\times$$
$$\times\ \varepsilon_{nn'n''n'''}\varepsilon_{rr'r''r'''}\varepsilon_{ss's''s'''}\bar{\varphi}\sigma_\mu^+\tau_m\varphi\bar{\varphi}'\sigma_\nu^+\tau_n\varphi'\times$$
$$\times\ \bar{\varphi}''\sigma_\rho^+\tau_r\varphi''\bar{\varphi}'''\sigma_\sigma^+\tau_s\varphi'''\bar{\varphi}\sigma_{\mu'}^+\tau_{m'}\varphi\bar{\varphi}'\sigma_{\nu'}^+\times$$
$$\times\ \tau_{n'}\varphi'\bar{\varphi}''\sigma_{\rho'}^+\tau_{r'}\varphi''\bar{\varphi}'''\sigma_{\sigma'}^+\tau_{s'}\varphi'''\bar{\varphi}\sigma_{\mu''}^+\tau_{m''}\varphi\bar{\varphi}'\times$$
$$\times\ \sigma_{\nu''}^+\tau_{n''}\varphi'\bar{\varphi}''\sigma_{\rho''}^+\tau_{r''}\varphi''\bar{\varphi}'''\sigma_{\sigma''}^+\tau_{s''}\varphi'''\bar{\varphi}\sigma_{\mu'''}^+\times$$
$$\times\ \tau_{m'''}\varphi\bar{\varphi}'\sigma_{\nu'''}\tau_{n'''}\varphi'\bar{\varphi}''\sigma_{\rho'''}^+\tau_{r'''}\varphi''\bar{\varphi}'''\sigma_{\sigma'''}^+\tau_{s'''}\varphi''' \tag{42}$$

(here, $\varepsilon_{\mu\nu\rho\sigma}$ is a completely antisymmetric tensor). Spinors χ_k^{α} are obtained from $\mathbf{O}$ by throwing off one $\bar{\varphi}_{\dot{\alpha} k}$, i.e.,

$$\chi_k^{\dot{\alpha}} = \frac{1}{4}\frac{\partial}{\partial\bar{\varphi}_{\dot{\alpha} k}}\mathbf{O} = \bar{\Phi}_k^{\dot{\alpha}}. \tag{43}$$

Spinors $\chi_k^{\dot{\alpha}}$ have $S = Y = 1$ and $i = 1/2$. They have the same transformation properties like spinors $\bar{\varphi}_{\dot{\alpha} k}$ (having strangeness $S = -1$) to which φ_k may already stick. As a result, the skeletons of so-called charmed particles arise:

$$O^{\Lambda_c^+} = \psi_{\alpha k}\otimes\chi_k,\quad O^{\Sigma_c} = (\vec{\tau}\psi_\alpha)_k\otimes\chi_k,$$

$$O^{\Xi_c} = (\gamma_\mu\psi)_{\alpha p}\bar{\varphi}_p\bar{\psi}\gamma_\mu\psi_k\otimes\chi_k,\quad O^{D^*} = \bar{\psi}\gamma_\mu\psi_k\otimes\chi_k,$$

$$O^{F^*} = \varphi_k\bar{\psi}_k\gamma_\mu\psi_m\otimes\chi_m,\quad O^{J/\psi} = \bar{\chi}_k\otimes\bar{\psi}_k\gamma_\mu\psi_m\otimes\chi_m. \tag{44}$$

We can give the following definition: skeleton O^Σ has charm $C = 1$ if it contains one 'heavy“ spinor $\chi_k^{\dot{\alpha}}$ or one additional variable χ_k. So we may call $\chi_k^{\dot{\alpha}}$ a charmed spinor (the usual spinor $\varphi_{\alpha k}$ is an uncharmed one).

Further, the masses of charmed particles calculated with (36) from [2] are

$$M_{\Lambda_c} = 2.48\ (2.28);\ M_{\Sigma_c} = 2.49\ (2.45);$$

$$M_{D^*} = 2.02\ (2.01);\ M_{\Xi_c} = 2.55\ (2.47);$$

$$M_{F^*} = 2.06\ (2.11);\ M_{J/\psi} = 2.88\ (3.1) \tag{45}$$

(in the brackets, experimental values are written). Hereat for these particles, we take the values of isotonic quantum number N equal to $\Lambda_c^+ : 32$, $\Sigma_c : 32$, $\Xi_c : 35$, $D^* : 33$, $F^* : 34$, $J/\psi : 64$, following from the explicit form of their skeletons O^Σ (44) (we recall that N is the total number of symbols $\varphi, \bar{\varphi}$ entering into O^Σ, see the definition of N in [2]).

Heavy quanta of strangeness χ_k (like light ones φ_k) may transit from the skeletons of fundamental particles to the skeletons of intermediate mesons η, π (comp. with Section 4). If χ_k is joined with $\eta, \vec{\pi}$, so a doublet D of mesons arises. If $\chi_k\varphi_k$ (or $\chi\vec{\tau}\varphi$) is joined with η (or $\vec{\pi}$), a singlet F arises.

Analogously, the quantum number of 'beauty“ B is connected with more heavy aggregates.

Conclusion

We consider that, without understanding the adequate reason of particle existence with the observed mass spectrum and kinds of interactions, it is impossible to plan experiments in the high-energy region. Without all that, high-energy physics looks like the electromagnetic theory of an electron which tried (before the creation of atomic physics and quantum theory) to explain the electron structure. In our opinion, without the explanation of atomic spectra concentrated in a narrow region of frequencies, it is

ISSN 0503-1265. Ukr. J. Phys. 2002. V. 47, N 7

impossible (using the mathematical apparatus of classical physics only) to build the adequate theory of electrons. The same, without the explanation of the nature of the particle spectrum (mainly concentrated in the limited energy region < 10 GeV), simply stepped over this region, to do any theoretical investigations in the physics of supersmall distances is impossible.

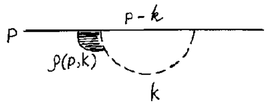
However, it is interesting to notice that, getting fundamental particles with all their characteristics (in particular, bilocality), their fundamental properties (the main of which is the relativity of motion in space-time), and kinds of interactions, we can deal with them only and forget for a certain time about the existence of the bi-Hamiltonian dynamical system from which they were originated. This fact testifies to a peculiar closedness of our particle world (from which the relativity principle originates) and renouncedness (isolatedness, separatedness) from it of such a physical reality as ether immediately described by the bi-Hamiltonian dynamics (so-called alpha-system) and its non-observability under usual conditions. Having three fundamental dimensional characteristics of ether, c, h, and k, we can, of course, obtain all particle characteristics [5]. But it may be said that to be occupied now in particle world is not already interesting. In our opinion, we have to plan the *experimentum crucis* in the high energy accelerator physics [2].

If this idea will be confirmed, our world outlook will become more complete and we will else have to do the last step towards the mathematical description of idos (embryo) of the Universe, the name of which is the Great Nothing, Creator of the *a priori* ether, or God.

APPENDIX

Here, we illustrate the calculation technique in the suggested strong interaction theory.

a) Renormalization of baryon mass. In the second order of perturbation theory, this process is described by the Feynman diagram



where the strong interaction formfactor $\rho(p, k)$ stands in that vertex from which a meson is emitted. The operator $\Sigma(p)$ of baryon proper energy is written in the form

$$\Sigma(p) = \frac{f^2}{k^2 (2\pi)^4 i} PP \int \frac{\hat{k}\gamma_5 (\hat{p} - \hat{k} + m)\hat{k}\gamma_5}{(k^2 - \mu^2)((p-k)^2 - m^2)} \times$$

$$\times\ \rho(p, k)\, d^4 k = A(p^2)\hat{p} + B(p^2)\, m.$$

[5] In particular, to calculate the exact masses of all elementary particles and the coupling constants (it seems to be a quite difficult problem like the calculation of the energy levels of a many-electron atom).

The $\Sigma(p)$ will be real if we use symmetric Green functions (see Section 3). As $\hat{k}\gamma_5(\hat{p} - \hat{k} + m)\hat{k}\gamma_5 = -k^2(\hat{p} + m) - (k^2 - 2pk)\hat{k}$, so under the simplifying condition $p^2 = m^2$ and neglecting the meson mass ($\mu^2 = 0$), we get the integrand $-\left(\frac{\hat{p} + m}{k^2 - 2pk} + \frac{\hat{k}}{k^2}\right)\rho(p, k)$. As $\rho(p, k) = \rho(p, -k)$ (see (46)), we obtain for A and B

$$A = B = \frac{i f^2}{k^2 (2\pi)^4} P \int \frac{\rho(p, k)}{n^2 - 2pk} d^4 k.$$

Representing the formfactor as (further, we use the system in which $c = h = k = 1$).

$$\rho(p, k) = \frac{1}{(2\pi)^2} \int e^{iqk}\, \delta(Q^2 + p^2)\, \delta(Q p) \times$$

$$\times\ \delta(q^2 + Q^2)\, \delta(q Q)\, d^4 Q d^4 q \qquad (46)$$

(here, the equality $\int e^{iqQ}\, \delta(q^2 + k^2)\, \delta(qk)\, d^4 q = = \int e^{iqk}\, \delta(q^2 + Q^2)\, \delta(q Q)\, d^4 q$ is used [1, 7]) and using the representation [9]

$$\frac{P}{A} = \frac{1}{2i} \int_{-\infty}^{\infty} \varepsilon(\alpha)\, e^{i\alpha A}\, d\alpha$$

for the symmetric propagator after the taking the Gauss quadratures over k_μ, we get

$$B = \frac{f^2}{32\, i\, \pi^2} \int_{-\infty}^{\infty} \frac{\varepsilon(\alpha)}{\alpha^2}\, d\alpha\, e^{-i\frac{(q - 2\alpha p)^2}{4\alpha}} \times$$

$$\times\ \delta(Q^2 + p^2)\, \delta(Q p)\, \delta(q^2 + Q^2)\, \delta(qQ)\, \frac{d^4 Q\, d^4 q}{(2\pi)^2}.$$

In view of the relations

$$\int e^{ipq}\, \delta(Q^2 + p^2)\, \delta(Q p)\, \delta(q^2 + Q^2)\, \delta(q Q) \times$$

$$\times\ \frac{d^4 Q\, d^4 q}{(2\pi)^2} = \frac{\sin p^2}{p^2}$$

(see [1]) and $p^2 = m^2$, we obtain

$$B = -\frac{f^2}{16\pi^2} \frac{\sin m^2}{m^2} \int_0^{\infty} \frac{d\alpha}{\alpha^2} \sin\left(\alpha m^2 + \frac{m^2}{4\alpha}\right) = \frac{f^2}{8\pi} N_1(m^2) \frac{\sin m^2}{m^2}, \qquad (47)$$

where N_1 is the Neumann function [13].

If we use the Feynman causal Green function and its Fock representation

$$\frac{1}{A} = -i \int_0^{\infty} e^{i\alpha A}\, d\alpha$$

we will obtain complex-valued result $i \frac{f^2}{8\pi} H_1^{(2)}(m^2) \frac{\sin m^2}{m^2}$ instead of a real one, where $H_1^{(2)}$ is the Hankel function of the second kind [13].

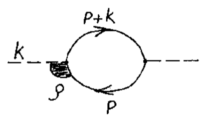
Complexity means that the Wick's S-matrix in non-unitary. Even if the Lagrangian is Hermitian (i.e., instead of the form-factor $\rho(p, k)$, we use the form-factor $\frac{1}{2}(\rho(p, k) + \rho(p - k, k)))$, the result will be still complex-valued: it exclusively depends on the form of strong formfactor. The solution of this trouble is in the following: to take the symmetric Green functions Δ^s as a particle propagator (see Section 3).

If $m^2 < 2.197$, $N_1(m^2) < 0$ (see [14]), and hence $B < 0$ (under this condition, $\frac{\sin m^2}{m^2} > 0$). This is very important as the formula for 'bare" baryon mass gives a too large value for a nucleon (see the mass table in [2]).

The renormalized fermion Green function is $\tilde{S}^{-1}(p) = S^{-1}(p) - \Sigma(p)$, where $S^{-1}(p) = \hat{p} - m$ is the non-renormalized function (see [9]). As $\Sigma(p) = B(\hat{p} + m)$ in the considered approximation, so $\tilde{S}^{-1}(p) = (1 - B)\hat{p} - (1 + B)m$. From here, we get $\tilde{m} = \frac{1 + B}{1 - B} m$ for renormalized mass and $\delta m = \tilde{m} - m \approx 2Bm$.

Applying this formula for a nucleon, we note that, according to the Section 4, the main contribution to the nucleon mass renormalization is given by η-meson (contribution from π-meson is 3/16 of that from η-meson, but contribution from K-meson is somewhat less). Non-renormalized nucleon mass is $m = 1.14$ (i.e., $m^2 = 1.3$), see [2]. As $N_1(1.3) = -0.55$, see [14], and $\frac{\sin 1.3}{1.3} = 0.74$ taking into account that $f^2/4\pi = 1/2$ in Heaviside's unit system, we get $B = -0.102$. Hence $\tilde{m} = 0.815m = 0.929\,\frac{kh}{c}$ ($khc = 1$ GeV). It is not far from the experimental value 0.938 $\frac{\text{GeV}}{c^2}$ [11].

b) M e s o n m a s s e s. In the suggested scheme, the masses of ' bare" gauge mesons are zero. These particles acquire masses as a result of dissociation into baryon-antibaryon pairs. In the second order of perturbation theory, this process is repsesented by the Feynman diagram



and is described by the expression

$$\Pi(k) = \frac{i f^2 I}{(2\pi)^4} PP \times$$

$$\times \int \frac{\mathrm{Sp}\,\hat{k}\gamma_5(\hat{p} + m)\hat{k}\gamma_5(\hat{p} + \hat{k} + m)}{(p^2 - m^2)((p + k)^2 - m^2)} \rho(p, k)\, d^4 p$$

(here the Fermi statistics of baryons is taken into account; I is the isotopic factor, see further). Representing the formfactor as

$$\rho(p, k) = \frac{1}{(2\pi)^2} \int e^{iQp}\, \delta(Q^2 + p^2)\, \delta(Qp) \times$$

$$\times\, \delta(q^2 + k^2)\, \delta(qk)\, d^4 Q\, d^4 q$$

and carrying out the integration over p, Q, q, we come to the following expression:

$$\Pi(k) = -\frac{f^2 I}{(2\pi)^2} \int_{-\infty}^{\infty} \frac{\varepsilon(\alpha)\varepsilon(\beta)}{(\alpha + \beta)^2} d\alpha\, d\beta \times$$

$$\times \exp\left[-i(\alpha + \beta)m^2 + i\frac{k^2\alpha\beta}{\alpha + \beta} - i\frac{k^2}{4(\alpha + \beta)}\right] \times$$

$$\times \left[-k^2 m^2 - i\frac{k^2}{\alpha + \beta} - (k^2)^2 \frac{\alpha\beta}{(\alpha + \beta)^2} - \frac{(k^2)^2}{4(\alpha + \beta)^2} + \right.$$

$$\left. + i\frac{(k^2)^2(\beta - \alpha)}{2(\alpha + \beta)^2}\frac{\partial}{\partial Z} - \frac{(k^2)^2}{2(\alpha + \beta)^2}\frac{\partial^2}{\partial Z^2}\right] \frac{\sin Z}{Z},$$

where $Z = k^2 \frac{\beta}{\alpha + \beta}$.

Product of two symmetric Green functions $S^s(X) S^s(X)$ may be represented as $\frac{1}{4}(S^c S^c + S^a S^a + S^c S^a + S^a S^c)$. Denoting the contribution in Π from $\frac{1}{4}(S^c S^c + S^a S^a) = \frac{1}{4}(S^+ S^+ + S^- S^-)$ (here, S^+, S^- are the positive and negative frequency parts of the perturbation function S, i.e., $\alpha > 0$ and $\beta > 0$ or $\alpha < 0$ and $\beta < 0$) as Π_1, we come to the following expression:

$$\Pi_1(k) = -\frac{f^2 I}{8\pi^2}\int_0^1 dz \int_0^\infty \frac{d\sigma}{\sigma}\left\{\cos\left(\sigma + \frac{A^2}{4\sigma}\right) \times\right.$$

$$\times \left[-k^2(m^2 + k^2 z(1 - z)) - \frac{A^4}{4\sigma^2}\left(1 + 2\frac{\partial^2}{\partial Z^2}\right)\right] +$$

$$\left. + \sin\left(\sigma + \frac{A^2}{4\sigma}\right)\left[-\frac{A^2}{\sigma} - \frac{A^2}{2\sigma}k^2(1 - 2z)\frac{\partial}{\partial Z}\right]\right\}\frac{\sin Z}{Z},$$

where $A = \sqrt{k^2(m^2 - k^2 z(1 - z))}$ and $Z = k^2 z$. Integration over σ and the use of the recurrent relations between Neumann functions give

$$\Pi_1(k) = -\frac{f^2 I}{8\pi^2}\int_0^1 dz\left[2\pi N_0(A)\left(k^2 m^2 + A^2\frac{\partial^2}{\partial Z^2}\right) + \right.$$

$$\left. + \pi A N_1(A)\left(k^2(1 - 2z) - 4\frac{\partial}{\partial Z}\right)\frac{\partial}{\partial Z}\right]\frac{\sin Z}{Z}.$$

At $k^2 \to 0$, we have

$$\Pi_1(k) = -\frac{f^2 I}{3\pi^2}\left[1 - k^2\frac{3m^2}{2}\ln\frac{\sqrt{k^2 m^2}\, e^{c - 1/6}}{2}\right]$$

where $C = 0.57$ is the Euler constant.

The contribution from $\frac{1}{4}(S^c S^a + S^a S^c) = \frac{1}{4}(S^+ S^- + S^- S^+)$ (i.e. when $\alpha > 0$ and $\beta < 0$, or $\alpha < 0$ and $\beta > 0$) is zero because S^+ and S^- as if extinguish each other.

It is very important to notice that as strong interaction is switched on just after the transition $f \to f^{\cdot}$, i.e. at very small distances (in the fiber) when only L-fields are arisen (R-fields have no time to arise), only

2-component spinor fields φ give the contribution in Π (but not 4-component bispinors $\begin{pmatrix}\varphi\\ \chi\end{pmatrix}$). Therefore, the previous result should be decreased by 2 times, and we have indeed,

$$\Pi(k) = \frac{f^2 I}{6\pi^2}\left[1 - k^2\frac{3m^2}{2}\ln\frac{\sqrt{k^2 m^2}\, e^{c-1/6}}{2}\right]. \qquad (48)$$

The renormalized propagator $\tilde{\Delta}(k)$ of a pseudoscalar particle is expressed in terms of the magnitude Π by the formula (see [9])

$$\tilde{\Delta}(k) = \Delta(k) - \Delta(k)\,\Pi(k)\,\Delta(k) + \ldots = \frac{\Delta(k)}{1+\Pi(k)\,\Delta(k)},$$

where $\Delta(k) = -\frac{1}{k^2}$ is a non-renormalized propagator. It follows from here that

$$\tilde{\Delta}(k) = \frac{-1}{k^2 - \Pi(k)} = \frac{-1}{Z_3(k^2)}\frac{1}{k^2-\mu^2}.$$

Here,

$$Z_3(k^2) = 1 + \frac{f^2 I m^2}{4\pi^2}\ln\frac{\sqrt{k^2 m^2}\, e^{c-1/6}}{2} \qquad (49)$$

is the renormalization constant of the propagator of a pseudoscalar field and

$$\mu^2 = \frac{f^2 I}{6\pi^2 Z_3(\mu^2)} \qquad (50)$$

is the mass square of the meson described by this field.

It has to note that the expression for Π is obtained in the second order of perturbation theory in the constant f. However, it contains some members of the fourth order in f. To take into account all members of the fourth order in f, we have to consider in (47), (48) m and f to be renormalized parameters, i.e. we have to substitute $m \to \tilde{m} = 0.94$ and $f \to \tilde{f}$ (constant f stays meanwhile nonrenormalized). As a result, we have the following numerical value for Z_3:

$$Z_3 = 1 + \frac{I\tilde{m}^2}{2\pi}\ln\frac{\sqrt{\tilde{m}^2 I/3\pi}\, e^{c-1/6}}{2}$$

(in Z_3, we put k^2 to be equal $\frac{f^2 I}{6\pi^2} = \frac{I}{3\pi}$ to find an approximate solution of the equation $k^2 = I/3\pi Z_3(k^2)$, hereby $f^2/4\pi = 1/2$). Therefore, $\mu^2 = I/3\pi Z_3$.

In the case of η and π mesons, the isotopic factor is obviously $I = 2$. For η-meson, we have

$$Z_3 = 1 - 0.31 = 0.69, \quad \mu^2 = \frac{2}{3\pi Z_3} = 0.31,$$

hence $\mu = 0.558\,\frac{kh}{c}$. The experimental value of the η-meson mass is 0.549 GeV/c^2 [11] ($khc = 1$ GeV).

From (35), it follows that we approximately have $m_\pi/m_\eta = \frac{1}{4} = 0.25$. The experimental value of this ratio is $\frac{135}{549} = 0.24$ [11].

The isotopic factor for K meson is $\left(\frac{1}{8}\right)^2 + 3\left(\frac{5}{8}\right)^2 = 1.2$ as follows from Section 4. Considering $m_\Lambda \approx m_\Sigma \approx m_N$, we obtain that, in this approximation, the ratio m_K/m_η is determined exclusively by the ratio of isotopic factors, i.e.

$$\frac{m_K}{m_\eta} = \sqrt{\frac{(1/8)^2 + 3\,(5/8)^2}{2}} \approx 0.8.$$

The experimental value of this ratio is $\frac{497}{549} = 0.9$ [11].

c) We stay else at the process of elastic scattering of baryons in the high energy region at large scattering angles. In the second order of perturbation theory for the cross section of this process, we have $d\sigma/dt \sim \frac{1}{s^6}\rho^2(p, q)$, where ρ is the strong formfactor (here, $s = (p_1 + p_2)^2$ is the total energy of two colliding particles in the center of mass system, $t = (p_1 - p'_1)^2 = q^2$ is the transferred momentum). As in this region $\rho(p, q) \approx \frac{1}{(pq)^2} \sim \frac{1}{s^2}$ (see [1]), we have $d\sigma/dt \sim \frac{1}{s^{10}}$. It is the well-known experimental result [15].

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Received 28.09.01