THE PROBLEM OF INTERACTIONS IN A DYNAMICAL THEORY OF PARTICLES (GENERAL QUESTIONS). 1

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Proceeding from the dynamics of a relativistic bi-Hamiltonian system based on the Heisenberg algebra $h_{16}^{(*)}$ (see [1, 4]), possible kinds of interactions between fundamental particles are derived. Three kinds of interactions: strong, electromagnetic, and gravitational ones are connected with the degeneration of the ground state $\dot{f}_z$ of the system and are described by the degeneration group $I = (\mathrm{SU}(2) \otimes U(1))_i \times U_e(1) \times \dot{T}_{3,1}$. The invariance group E (hidden symmetry) of the state is defined. The space-time structure of interacting particle states (bilocal fields) is investigated and equations for these bilocal fields are obtained.

1. Introduction

The fundamental problem of the local interacting field theory is ultraviolet divergences. The hope of reducing these by postulating so-called supersymmetric partners was not justified: no signal of their existence was observed. Another generalization of the theory connected with the so-called string model met serious troubles upon quantization[1]. Concerning the well-known composite models of particles (quarks, preons, and so on), it can be said that the mentioned troubles and a new so-called confinement problem arise from these models.

At present it is considered that the particle constituents (namely quarks) may be kept together (within the particle) solely by means of an interaction. In connection with this, the colour quark-gluon interaction is considered to be fundamental (as is the weak one at the quark level). In other words, the particle constituents must interact in order to form the particle and its interactions. It turns out that this is not obligatory: there is another possibility for constituent confinement (see [1]). Moreover, this phenomenon is impossible at all in the framework of Lagrangian or Hamiltonian systems (quarks are considered as a Lagrangian field system). In this perspective, a new solution to the problem of ultraviolet divergences is suggested not transferring the issue but approaching it at that place where it arose.

[1]We consider that at the classical level elementary particle must be considered to be point like object (concerning quantum level see further).

Rejecting the notion of interacting constituents, a fundamentally another reason of particle existence is considered (see [1]). In the theory developed in [1], neither particle constituents (granules) nor their dynamical structures (quanta f and $\dot{f}$) interact: it is a level of so-called asymptotic freedom (to be referred as the subquantum level). At this very deep level, there are no interactions at all (all interactions except the gravitational one are switched off, see further and [2]). The interactions appear at first only at the fundamental particle level, i.e. after the transition $f \to \dot{f}$ which happens in the bi-Hamiltonian dynamical system, resulting in the appearance of fundamental (Lagrangian) fields ψ^{Σ} and interactions between them. Therefore we call them fundamental. The arisen inreactions are inherent to the Lagrangian particles only (and Lagrangian systems at all; granules are not a Lagrangian system because, at their level, the space-time structure is discontinuous).

In connection with this, it is important to emphasize that the direct reason of ultraviolet divergences is the use of non-proper (generalized) functions (like Dirac's delta-function) in the second quantization method, see [1]. The existence of such objects is connected with the concept of space-time as a continuum endowed by the Newtonian differentiable structure and the Lebesgue measure. It is to be admitted that the Newtonian understanding of space is only a mathematical model: it is insufficient as a description of physical space (see [1]) and as a basis for the consequent quantized field theory and the theory of elementary particles. Analysis of the structure of space-time at very small distances [1] leads to a concept of physical space-time based on a new kind of dynamics, which may underlie the fundamental particle theory and their interactions [1]. Here we are outgoing from this concept.

Obviously, speaking about fundamental interactions, it is needed to have the well-defined zero-approximation, i.e. the notion of 'bare“ (non-interacting) object, interactions of which are not postulated but derived from the nature of the object. The condition of being fundamental is that these interactions may not be reduceable to more fundamental ones, namely then the zero approximation is another. This program is realized in the framework of a new approach to the particle problem [1]. Hereby it is very important to emphasize that the bi-Hamiltonian dynamical system (algebra $h_{16}^{(*)}$ with SU(2) isotopic symmetry) does not contain quarks and gluons at all (as objects of SU(3) group), but nevertheless we may speak about the highest (SU(3), SU(4) ¯ and so on) approximate symmetries in the spectrum of fundamental particles. Here some general topics of the interaction problem are considered.

2. Discrete Symmetry, Degeneration Group, and Hidden Symmetry of the State $\dot{f}_z$

1) We begin our consideration from the discussion of some discrete symmetries of the ground state $\dot{f}_z$ of the bi-Hamiltonian dynamical system based (for simplicity) on the Heisenberg algebra $h_8^{(*)}$ [1]. We recall that the generators of the $h_8^{(*)}$-algebra obey the commutation relations

$$[\Phi_\alpha, \overline{\Phi}_\beta] = \delta_{\alpha\beta}, \quad [\Phi_\alpha, \Phi_\beta] = [\overline{\Phi}_\alpha, \overline{\Phi}_\beta] = 0 \qquad (1)$$

and are written in the form of [1]

$$\Phi = \begin{pmatrix} \varphi_\alpha \\ \partial/\partial\overline{\varphi}_\alpha \end{pmatrix}, \quad \overline{\Phi} = \left(-\frac{\partial}{\partial\varphi_\alpha}, \overline{\varphi}_\alpha\right), \quad \alpha = 1, 2. \qquad (2)$$

Like the Dirac's theory, we denote the discrete symmetries as P, C, T and define them on the generators $\Phi, \overline{\Phi}$ by the expressions

$$P: \ \Phi \to \Phi^P = \gamma_4 \Phi, \quad \overline{\Phi} \to \overline{\Phi}^P = \overline{\Phi}\gamma_4,$$

$$C: \ \Phi \to \Phi^C = C^{-1}\overline{\Phi}', \quad \overline{\Phi} \to \overline{\Phi}^C + C\Phi',$$

$$T: \ \Phi \to \Phi^T = i\gamma_4\gamma_5 C^{-1}\overline{\Phi}', \quad \overline{\Phi} \to \overline{\Phi}^T = i\gamma_4\gamma_5 C\Phi', \qquad (3)$$

where Φ' is a transposed magnitude. Here C is the so-called charge conjugation matrix obeying the conditions: $C' = -C, C^+ = C^{-1} = -C^*$, $C^2 = -1$. Its explicit form in the Weyl representation is $C = \begin{pmatrix} \varepsilon & 0 \\ 0 & \dot{\varepsilon} \end{pmatrix}$ where $\varepsilon = -\dot{\varepsilon} = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}$ is a 2×2-matrix [3]. In this representation, we have $\gamma_4 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}$, $\gamma_5 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}$ where 1 and 0 are the 2×2 unit and zero matrices. It follows from (3) that $PT: \Phi \to i\gamma_5 C^{-1}\Phi'$, $CP: \Phi \to \gamma_4 C^{-1}\overline{\Phi}'$, TC: $\Phi \to i\gamma_4\gamma_5\Phi$, and $CPT: \Phi \to i\gamma_5\Phi$. Of course, we may write

$$\Phi^P = T_P \Phi T_P^{-1}, \quad \Phi^C = T_C \Phi T_C^{-1},$$

$$\Phi^T = T_T \Phi T_T^{-1}, \qquad (4)$$

where T_P, T_C, T_T are the operators acting in the representation spaces $\mathbf{F}, \dot{\mathbf{F}}$ of the algebra $h_8^{(*)}$.

It follows from (1) that the algebra $h_8^{(*)}$ is invariant under the reflections P, C, T if we take the Schwinger's rule to read the commutation relations from the right to the left in the case of T-reflected generators, as they obey the relations $[\Phi_\alpha^T, \overline{\Phi}_\beta^T] = -\delta_{\alpha\beta}$. However the spaces $\mathbf{F}, \dot{\mathbf{F}}$ and belonging to them states of the relativistic bi-Hamiltonian system f and $\dot{f}$ are not invariant under these transformations. Indeed, it follows from (2) and (3) that

$$P: \varphi \to \varphi^P = \frac{\partial}{\partial\overline{\varphi}}, \quad C: \varphi \to \varphi^C = -\varepsilon\frac{\partial}{\partial\varphi},$$

$$T: \varphi \to \varphi^T = -i\dot{\varepsilon}\,\overline{\varphi}. \qquad (5)$$

We see that the C-transformation is connected with the complex Fourier transformation $\mathcal{F}$, i.e. $T_C = \mathcal{F}$. If $f(\varphi) \in \mathbf{F} = \mathcal{F}_\mathrm{F} \otimes \mathcal{F}_0$ (the structure of the carrier space $\mathbf{F}$ is determined in [1]), so

$$\mathcal{F}: f(\varphi) \to (T_C f)(\varphi') =$$

$$= \int e^{\varphi'\varepsilon\varphi + \overline{\varphi}'\dot{\varepsilon}\overline{\varphi}} f(\varphi)\, d\mu(\varphi), \qquad (6)$$

where the measure $d\mu(\varphi) = \prod_{\alpha=1,2} d\varphi_\alpha \wedge d\overline{\varphi}_\alpha$ (it is invariant under the Lorentz transformations $\varphi_\alpha \to \upsilon_\alpha^\beta \varphi_\beta$, $\upsilon \in \mathrm{SL}_l(2, \mathbf{C})$). Hence, the C-transformation differs essentially from its Dirac's analog. Obviously, the ground state of the system $\dot{f}_z = C e^{\overline{\varphi}z - \overline{z}\varphi}$ (see [1]), where $\varphi \doteq \varphi_2$ is an additional variable in the extended Fock representation (see [1]) is not invariant under the C and T (and hence P) transformations.

It is important to emphasize that T-assymetry of the additional variable φ (see (5)) leads to the degeneration (doubling) of this variable. It turns out that this leads, first, to the doubling of Universe (besides our Universe, there exists the antiUniverse, see [2]) and, secondly, to the consideration of the Heisenberg algebra with doubling dimension $h_{16}^{(*)}$, containing the real SU(2)-isotopic symmetry of our Universe, see [4]. Further we consider the realistic model of a bi-Hamiltonian system based on the $h_{16}^{(*)}$-algebra. With T-asymmetry of the suggested scheme, the fermion-antifermion asymmetry of our Universe or so-called selection rules are connected [4] (selection rule is: in our Universe, all additional variables $\overline{\varphi}_k$ must be pairing with $\varphi_{\alpha k}$ or φ_k; it means that only fermions may arise in the transition $f \to \dot{f}$). We will see further that P and C symmetries will be restored at the level of massive fundamental (Lagrangian) particles.

2) Now we go over to the continuum degeneration of the state $\dot{f}_z$. We define the generation group $\dot{I}$ as a stationary group of the 4-momentum $\dot{p}_\mu = -i\overline{\Phi}\gamma_\mu P_- \Phi$ of the ground state $\dot{f}_z$. First of all, it consists of the $\dot{T}_{3,1}$-transformations (generators $\dot{p}_\mu$), the small Lorentz group $(\mathrm{SL}_l(2, \mathbf{C}) \otimes U_l(1))_{\dot{p}}$ of $\dot{p}_\mu$ (it is subgroup of the group $\mathrm{GL}_l(2, \mathbf{C})$, generators of which are $I_{\mu\nu} = \overline{\Phi}\Sigma_{\mu\nu}\Phi$, $A = \overline{\Phi}\Phi$, $B = \overline{\Phi}\gamma_5\Phi$), further the isotopic group $(\mathrm{SU}(2) \otimes U(1))_i$ (generators $\vec{i} = -\frac{1}{2}\overline{\Phi}\vec{\tau}\Phi$, $Y = -\overline{\Phi}\Phi - 2$), and the electromagnetic group $U_e(1)$ (generator $Q = \frac{1}{2}Y + i_3 = -\overline{\Phi}q\Phi - 2$, where $q = \frac{1}{2}(1 + \tau_3) = \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix}$). The latter is a subgroup of $U(2) = \mathrm{SU}(2) \otimes \overline{U}(1)$ (it should not be mixed with the isotopic group $(\mathrm{SU}(2) \otimes U(1))_i$). All these generators commute with $\dot{p}_\mu$, so the transformations connected with them do not change the 4-momentum $\dot{p}_\mu$ of the ground state $\dot{f}_z$ (we emphasize that the dilatation $H_i(1)$ is not included into the degeneration group $\dot{I}$; also $H_l(1)$ is not included in I). Under the action of these transfrormations, the uncountable number of states $\{e^{i\rho a}\dot{f}_{u^+z}\}$ with the same 4-momentum $\dot{p}_\mu$ arise. Here, $u = \exp(\frac{i}{2}\vec{\tau}\vec{\theta} + i\theta)\, e^{i\chi q}$ and $a_\mu, \theta, \vec{\theta}, \chi$ are parameters of the degeneration group $\dot{I} = \dot{T}_{3,1} \otimes (\mathrm{SU}(2) \otimes U(1))_i \times U_e(1)$ (χ is the phase of the magnitude $\varepsilon = \frac{z_1}{z_2}$, see further). Therefore the degenerate states obey the same equation that $\dot{f}_z$. All these parameters are open in the theory (unlike the parameters of the group of hidden symmetry, see below). They are elastic and may be exited. With them are connected the interactions between particles (see further).

3) We determine the group of the hidden (exact) symmetry E of a realistic bi-Hamiltonian system as the invariance group of the ground state $\dot{f}_z = C\, e^{\bar{z}_k\varphi_k - \overline{\varphi}_k z_k}$. As the additional variables $\varphi_k, \overline{\varphi}_k$, which the state $\dot{f}_z$ depends on, are scalars of the $\mathrm{GL}_l(2\,\mathbf{C})$-group (see [1]), so E includes the latter. E includes also the stationary group W_z of the isospinor $z = \begin{pmatrix} z_1 \\ z_2 \end{pmatrix} : w_z z = z,\ w_z \in W_z$. Therefore we have $T(w_z)\dot{f}_z = \dot{f}_z$. So we have $E = \mathrm{GL}_l(2, \mathbf{C}) \otimes W_z$. It is important to note that transformations from the group E spread the eigenvalue ρ_μ of the 4-momentum $\dot{p}_\mu$ on the lower area of the light cone N_- (see [1]).

It is not difficult to verify that W_z is the two-parametric group of matrices $w_z = \begin{pmatrix} a & \varepsilon(1-a) \\ c & 1-\varepsilon c \end{pmatrix}$, where $\varepsilon = \frac{z_1}{z_2}$. It is a subgroup of the group $\mathrm{GL}(2, \mathbf{C})$ isomorphic to the group of affine transformations of the straight line with generators $q_1 = \begin{pmatrix} \varepsilon & -\varepsilon^2 \\ 1 & -\varepsilon \end{pmatrix}$, $q_2 = \begin{pmatrix} 1 & -\varepsilon \\ 0 & 0 \end{pmatrix}$ obeying the relations $[q_1, q_2] = q_1$.

The intersection $I \cap E = (\mathrm{SL}_l(2, \mathbf{C}) \otimes U_l(1))_{\dot{p}} \otimes U_\varepsilon$ plays on important role in the theory. It is the symmetry group of the state $\dot{f}_z(\dot{x}) = e^{i\rho\dot{x}}\dot{f}_z$. It is interesting to note that the intersections $\mathrm{SU}_i(2) \cap E$ and $U_i(1) \cap E$ consist of only the unit element, but the intersection $(\mathrm{SU}(2) \otimes \overline{U}(1)) \cap E = U_\varepsilon$, where U_ε is the group of matrices $e^{iq_\varepsilon\alpha}$ and $q_\varepsilon = \frac{1}{\sqrt{1+|\varepsilon|^2}} \times \begin{pmatrix} 1 & -\varepsilon \\ -\varepsilon & |\varepsilon|^2 \end{pmatrix}$, so that $q_\varepsilon z = 0$ and $q_\varepsilon^+ = q_\varepsilon$, $q_\varepsilon^2 = q_\varepsilon$, $\mathrm{Tr}\, q_\varepsilon = 1$. Hence we have $u_\varepsilon = 1 + (d-1)\, q_\varepsilon$, where $d = \det u_\varepsilon = e^{i\alpha}$. The group U_ε is a subgroup of the group $U(2) = \mathrm{SU}(2) \otimes \overline{U}(1)$. Parameters of the group $I \cap E$ do not appear in the theory: they are hidden at any level. Therefore they may not be exited and there are no interactions connected with them. In particular, there is no so-called torsion fields (parameters of the $\mathrm{SU}_l(2)$ group) connected with spin (comp. with [5]) and gauge weak interaction connected with the parameters of the group $U(2)$ (comp. with [6]), with the exception of the parameter χ of the subgroup $U_e(1)$ of the group $U(2)$ (concerning a new theory of weak interaction, see [7]).

4) the 4-momentum p_μ of the states f^Σ has a quite analogous degeneration group. In this case, degeneration is connected with the harmonic analysis on the space $\mathbf{F}$ and leads to the forming of finite-dimensional multiplets $\Sigma = (s, F; i, Y, Q)$ of the group $I = (\mathrm{SL}_l(2, \mathbf{C}) \otimes U_l(1))_p \otimes (\mathrm{SU}(2) \otimes U(1))_i$. With the degeneration group $T_{3,1}$, the switching on of the gravitational interaction between quanta f is connected.

3. Restoring Some Symmetries at the Fundamental Particle Level

1) Due to the 100 percent fermion-antifermion asymmetry of the suggested scheme (see [4]), only fermions arise in the quantum transition $f \to \dot{f}$

(antifermions do not arise). Hereby there appear fields with positive frequency and L-helicity only. Such fields are written in the form

$$\overset{(+)}{\psi}{}^{\Sigma}_{\alpha}(X, Y) = \frac{1}{(2\pi)^{3/2}}\int e^{iPX}\,\theta(P_0)\,\delta(P^2 + M^2_{\Sigma})\times$$

$$\times\, a^{\Sigma}_{\alpha}(P, Y)\, d^4P, \tag{7}$$

where

$$\theta(P_0)\, a^{\Sigma}_{\alpha}(P, Y) = \frac{1}{2\pi}\int e^{iQY}\,\theta(P_0 + Q_0)\,\theta(P_0 - Q_0)\times$$

$$\times\,\theta(-P^2)\,\delta(Q^2 + P^2)\,\delta(QP)\, O^{\Sigma}_{\alpha}(P, Y)\, d^4Q \tag{8}$$

(see [1]). Looking at these expressions, we can say that the physical space-time is a fibration of the base $\mathbf{A}_{3,1}\supset X$ (Poincare affine space) with a fiber $\mathbf{R}_{3,1}\otimes S$ where $\mathbf{R}_{3,1}\supset Y$ is the Minkowski vector space. Here $S = \underset{\Sigma}{\oplus} S^{\Sigma}$ is the Whitney sum. Elements of the fiber are denoted as $\psi_{\alpha, Y}$. The sections of this fibration are the fields $\psi_{\alpha, Y}(X)$ which we prefer to denote as $\psi_{\alpha}(X, Y)$ and call bilocal (it has to note that they should not be mixed up with the Yukawa bilocal fields [8]). (Sometimes we will call such a quantum object as the Lorentzian membrane; representing particles by Feynman's diagrams, we can say that a particle is an enough complicated 'cable" but not a simple Feynman's 'line".) We can say that the size of the space $\mathbf{R}_{3,1}$ is effectively limited by the length k^{-1}, see [1] (hence the fundamental length exists there in the inner space but not in the external one).

In (7), $a^{\Sigma}_{\alpha}(P, Y)$ is a field amplitude normalized in usual manner (one particle in the whole space). As is well known, the possibility of arbitrary normalization of a wave function ψ deprives it of the immediately physical sense and leads to the probability interpretation of ψ to be a wave of probability or information. It turns out this fundamental property is caused by the non-unitary character of the bi-Hamiltonian theory, namely by the non-Hermitian form $\langle \dot{f}_z, f^{\Sigma}\rangle$ defining the transition matrix element: in a geometry with non-Hermitian form there is no notion of the vector length [9]. Another property – reduction – is connected with the interaction and measurement problem.

It follows from the explicit form of $\overset{(+)}{\psi}{}^{\Sigma}_{\alpha}(X, Y)$ that the field $\overset{(+)}{\psi}{}^{\Sigma}_{\alpha}(X, Y)$ obeys the Klein – Gordon equation $(\Box_X - M^2_{\Sigma})\,\psi^{\Sigma} = 0$. However not this but the equations of first order are evolution equations. It will be seen better in the case of the Dirac field $\overset{(+)}{\psi}(X, Y)$ of L-helicity. As is well known, there is no Lagrangian for this field. For the construction of the Lagrangian, some supplementary fields are needed, namely R-fields, which are given by the Dirac equations $i\,\overset{+}{\sigma}_{\mu}\frac{\partial}{\partial X_{\mu}}\overset{(+)}{\varphi} = -m\,\overset{(+)}{\chi}$. So the R-field χ appears as a result of evolution in space-time of the L-field φ. Together φ and χ give the bispinor Dirac field $\overset{(+)}{\psi}(X, Y) = \begin{pmatrix}\overset{(+)}{\varphi}(X, Y)\\ \overset{(+)}{\chi}(X, Y)\end{pmatrix}$. It has positive frequency like $\overset{(+)}{\varphi}$.

2) Under complex conjugation (involution), the field $\overset{(+)}{\psi}$ transits into the negative frequency field $\overset{(-)}{\psi}$. In the quantized field theory, $\overset{(+)}{\psi}$ and $\overset{(-)}{\overline{\psi}}$ describe the particle creation and antiparticle annihilation processes. We emphasize that the negative frequency part $\overset{(-)}{\psi}$ is not appeared in the transition $f \to \dot{f}$, it may appear only as a result of an interaction (coupling) of the field $\overset{(+)}{\psi}$ with the field mediating the interaction when the latter takes the energy of the field $\overset{(+)}{\psi}$ and it transits into the state with negative energy, i.e. as a result of the quantum transition taking a place in the Lagrangian field system. In the quantized field theory, $\overset{(-)}{\psi}$ and its complex conjugation $\overset{(+)}{\overline{\psi}}$ describe the processes of annihilation of a particle and creation of an antiparticle [10]. The complete field is the sum

$$\psi^{\Sigma} = \overset{(+)}{\psi}{}^{\Sigma} + \overset{(-)}{\psi}{}^{\Sigma}, \quad \overline{\psi}^{\Sigma} = \overset{(+)}{\overline{\psi}}{}^{\Sigma} + \overset{(-)}{\overline{\psi}}{}^{\Sigma} \tag{9}$$

hereby

$$\overline{\overset{(+)}{\psi}{}^{\Sigma}} = \overset{(-)}{\overline{\psi}}{}^{\Sigma}, \quad \overline{\overset{(+)}{\overline{\psi}}{}^{\Sigma}} = \overset{(-)}{\psi}{}^{\Sigma} \tag{10}$$

(it is important to note that $\overset{(-)}{\overline{\psi}}{}^{\Sigma} \neq \overset{(+)}{\overline{\psi}}{}^{\Sigma}$). As a result, we obtain the following expressions:

$$\overline{\psi}^{\Sigma}_{\alpha k}(X, Y) = \frac{1}{(2\pi)^{3/2}}\int e^{iPX}\,\delta(P^2 + M^2_{\Sigma})\times$$

$$\times\,[\theta(P_0)\, a^{\Sigma}_{\alpha k}(P, Y) + \theta(-P_0)\,\overline{b}^{\Sigma}_{\alpha k}(-P, Y)]\, d^4P, \tag{11}$$

$$\overline{\psi}^{\Sigma}_{\alpha k}(X, Y) = \frac{1}{(2\pi)^{3/2}}\int e^{iPX}\,\delta(P^2 + M^2_{\Sigma})\times$$

$$\times\,[\theta(-P_0)\,\overline{a^{\Sigma}_{\alpha k}(-P, Y)} + \theta(P_0)\, b^{\Sigma}_{\alpha k}(P, Y)]\, d^4P, \tag{12}$$

which are used in the phenomenological approach to quantized field theory [10].

Although the R-field arises only as a result of evolution of the L-field and has no 'skeleton" (see [1]; we recall that the particle skeleton is a symbol $O^{\Sigma}(\varphi)$ constructed from the spinors $\varphi_{\alpha k}, \bar{\varphi}_{\alpha k}$ and additional variables $\varphi_k, \bar{\varphi}_k$ only), we may speak about the skeleton of the R-field which is built from the R-spinor

$$\chi_k^{\dot{\alpha}} = \frac{\overset{+}{\sigma}{}_{\mu}^{\alpha\beta} \rho_{\mu}}{\sqrt{2\pi\rho}} \varphi_{\beta k}, \qquad (13)$$

where π_μ and ρ_μ are the 4-momenta of the states f^{Σ} and $\dot{f}_z$, correspondingly (at $\rho_\mu = 0$, we obviously have $\chi = 0$). Namely such a spinor will be obtained from the Dirac equation written in the momentum representation $\overset{+}{\sigma}_\mu (\pi - \rho)_\mu \varphi(\pi) = -\overset{+}{\sigma}_\mu \rho_\mu \varphi(\pi)$ as $\overset{+}{\sigma}_\mu \pi_\mu \varphi(\pi) = 0$ because $\varphi(\pi) = \bar{\sigma}_\mu \pi_\mu a$. Here $a = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$ is a constant spinor, see [1]. Obviously, we may speak about the spinor χ (13) only after the transition $f \to \dot{f}$ (χ depends on ρ_μ).

On the bispinor $\psi = \begin{pmatrix} \varphi \\ \chi \end{pmatrix}$, one may define the P'-reflection by means of the transformation: $\psi \to \gamma_4 \psi$, where γ_4 is the Dirac matrix and thereby to restore the P-symmetry broken at the subquantum level (see above).

3) It is not difficult to show using bespinors $\psi = \begin{pmatrix} \varphi \\ \chi \end{pmatrix}$, where χ is given by the formula (13), and difinition of transition amplitudes [1] that $\bar{\psi}\gamma_\mu\gamma_5\tau_k\psi \sim Q_\mu \bar{z}\tau_k z$ and $\bar{\psi}\gamma_\mu\tau_k\psi \sim P_\mu \bar{z}\tau_k z$. As the integral

$$\int Q_\mu e^{iQY} \theta(P_0 + Q_0)\,\theta(P_0 - Q_0) \times$$
$$\times\, \delta(Q^2 + P^2)\,\delta(QP)\, d^4Q$$

is equal zero at $Y = 0$, so one can conclude that, in the asymptotic limit $|X| >> |Y|$ (i.e. at the large distances from the place of field creating), pseudovector fields are zero. It means that there are no pseudovector fundamental particles in the suggested scheme (that is in accordance with experimental data, see [11]). On the contrary, the vector fields

$$P_\mu \int e^{iQY} \theta(P_0 + Q_0)\,\theta(P_0 - Q_0) \times$$
$$\times\, \delta(Q^2 + P^2)\,\delta(QP)\, d^4Q$$

differ from zero at $Y = 0$.

It turns out that there are no pseudoscalar fundamental particles in the suggested scheme too. Indeed, the pseudoscalar combinations ($\overset{+}{\rho} = \overset{+}{\sigma}_\mu \rho_\mu$)

$$\bar{\psi}\gamma_5\tau_k\psi = \bar{\chi}\tau_k\varphi - \bar{\varphi}\tau_k\chi =$$

$$= \frac{1}{\sqrt{2\pi\rho}}(\bar{\varphi}\overset{+}{\rho}\tau_k\varphi - \bar{\varphi}\overset{+}{\rho}\tau_k\varphi) \equiv 0$$

are identically zero. This means that η and π (and also K) mesons may not be fundamental particles. For them, there is only one possibility to be quanta of degeneration fields (see below). However the scalar combinations

$$\bar{\psi}\tau_k\psi \sim \sqrt{2\pi\rho}\,\bar{z}\tau_k z \sim \sqrt{-P^2} = M_\Sigma$$

differ from zero if the mass of particle is not zero.

In general case (arbitrary Σ), the field $\overset{(+)}{\psi}{}^{\Sigma}(X, Y)$ has the skeleton $O^{\Sigma}(\varphi)$ which is transformed under a finite-dimensional irreducible representation $\left(\frac{n}{2}, \frac{m}{2}\right)$ of the group $SL_l(2, \mathbf{C})$. The skeleton $O^{\Sigma}(\varphi)$ is built from L-spinors φ (($\frac{1}{2}$, 0)-representation) and complex conjugate $\bar{\varphi}$ ((0, $\frac{1}{2}$)-representation). Using R-spinors (13), we can build all other fields connected with $\overset{(+)}{\psi}{}^{\Sigma}(X, Y)$ which give the possibility to build the total field $\overset{(+)}{\Psi}{}^{\Sigma}(X, Y)$ obeying the first-order differential equation [12]

$$\left(\Gamma_\mu^\Sigma \frac{\partial}{\partial X_\mu} + M_\Sigma\right)\Psi^\Sigma = D^\Sigma\left(\frac{\partial}{\partial X}\right)\Psi^\Sigma = 0. \qquad (14)$$

It is the true evolution equation of the field Ψ^Σ (but not the Klein – Gordon one). It is shown in [12] that the operator $D^\Sigma\left(\frac{\partial}{\partial X}\right)$ has such a divisor $d^\Sigma\left(\frac{\partial}{\partial X}\right)$ that $d^\Sigma D^\Sigma = \Box_X - M_\Sigma^2$ (the Klein – Gordon operator). Interactions between fundamental particles are switching on in the frame of the first-order equation (14). Now we go over to consideration of this question.

4. Switching on the Interactions

Here we give the general schematic description of the procedure of switching on the interactions without any specification of their kind (details will be given in a specific consideration of that or other kind of interaction).

1) All the collection of degenerate states $\dot{f}_z$ may be written as $\{T(g)\dot{f}_z\}$, where $g = e^{i\Phi}$ is an element of the degeneration group I hereby the phase $\Phi = \sum_i \theta_i q_i$, where θ_i are the parameters of the group I and q_i (matrices) are its generators (charges). Interactions are switched on when the parameters θ_i become any functions $\theta_i(X, Y)$ on a fibration $(\mathbf{A}_{3,1}, \mathbf{R}_{3,1})$ ($X \in \mathbf{A}_{3,1}$ is the base — the Poincare affine space-time, $Y \in \mathbf{R}_{3,1}$ is a fiber — the Minkowski vector space, see [1]). So I becomes a locally acting Lie group (comp. with the Yang — Mills mechanism [13]). We may say that interactions are a consequence of stimulation of degeneration channels (parameters) θ_i when the energy released in the quantum transition $f \to \dot{f}$ comes into these channels. As a result, the parameters θ_i being elastic will depend on the coordinates x_μ, $\dot{x}_\mu$ or X_μ and Y_μ (see [1]), i.e. become some fields. Hence the interactions we considered arise after the transition $f \to \dot{f}$ and produce the particle fields $\Psi^\Sigma(X, Y)$. This kind (gauge) of interactions is inherent in only the fundamental particles originated in the transition $f \to \dot{f}$.[2]

In the transition matrix element $\langle e^{i\hat{\Phi}}\dot{f}_z, f^\Sigma\rangle$, the operator $e^{i\hat{\Phi}}$, where $\hat{\Phi} = T(\Phi)$, may be transferred from $\dot{f}_z$ to f^Σ. As f^Σ are the eigenvectors of the operators (charges) $\hat{Q}_i = T(q_i) : \hat{Q}_i f^\Sigma = Q_i^\Sigma f^\Sigma$, so we have $\langle e^{i\hat{\Phi}}\dot{f}_z, f^\Sigma\rangle = e^{-i\Phi^\Sigma}\langle \dot{f}_z, f^\Sigma\rangle$, where $\Phi^\Sigma = \sum_i \theta_i(\dot{x}) Q_i^\Sigma$. Obviously, not only the matrix elements $\langle \dot{f}_z, f^\Sigma\rangle$ but the transition amplitudes $\langle\langle \dot{f}_z, f^\Sigma\rangle\rangle$ as well as the particle fields $\Psi^\Sigma(X, Y)$ undergo a phase modulation. So, at the switching on the interactions, particle fields are modulated according to the formula:

$$\Psi^\Sigma(X, Y) \to e^{-i\Phi^\Sigma(X,Y')}\Psi^\Sigma(X, Y) = \tilde{\Psi}^\Sigma(X, Y, Y').$$

Here the functions $\Phi^\Sigma(X, Y')$ play the role of connectness in the fibration $(\mathbf{A}_{3,1}, \mathbf{R}_{3,1})$, whose sections are the fields $\Psi^\Sigma(X, Y)$. From the evolution equations (14) for the 'bare" (non-interacting) fields $\Psi^\Sigma(X, Y)$, the equations for the interacting (modulated) fields $\tilde{\Psi}^\Sigma(X, Y, Y')$ follow. In a correct interaction theory, the fields $\theta_i(X, Y')$ must be considered as weak and slowly changing. Therefore the interaction equations are written in the form

$$\left(\Gamma_\mu^\Sigma \frac{\partial}{\partial X_\mu} + M_\Sigma\right)\Psi^\Sigma(X, Y) =$$

$$= -i\Gamma_\mu^\Sigma\left(\frac{\partial}{\partial X_\mu}\Phi^\Sigma(X, Y')\right)\Psi^\Sigma(X, Y). \tag{15}$$

Under the condition mentioned above, we may consider $\tilde{\Psi}^\Sigma(X, Y, Y')$ to be the 'bare" field $\Psi^\Sigma(X, Y)$ that we made already in (15).

2) In the case of $T_{3,1}$-degeneration of quanta f (here we do not consider $\dot{T}_{3,1}$-degeneration of $\dot{f}$), we have in the asymptotic limit $|X| >> |Y| : \frac{\partial}{\partial X_\mu}\Phi^\Sigma = \frac{\partial a_\nu}{\partial X_\mu}\frac{\partial}{\partial X_\nu}$. The Lagrangian of this interaction is written as $h_{\mu\nu}^{(0)} T_{\mu\nu}^\Sigma$, where $T_{\mu\nu}^\Sigma = \frac{i}{4}\left[\overline{\Psi}^\Sigma \Gamma_\mu^\Sigma \frac{\partial}{\partial X_\nu} - \left(\frac{\partial}{\partial X_\nu}\overline{\Psi}^\Sigma\right)\Gamma_\mu^\Sigma\right]\Psi^\Sigma$ is the energy-momentum tensor of the field Ψ^Σ and $h_{\mu\nu}^{(0)} = \frac{1}{2}\left(\frac{\partial a_\nu}{\partial X_\mu} + \frac{\partial a_\mu}{\partial X_\nu}\right)$ is the so-called gauge gravitational potential. The more general phase Φ depending on the path in $\mathbf{A}_{3,1}$ corresponds to the real gravitational potential. Concerning the definition of the Newtonian gravitational constant γ, see [14].

Quite analogously under the $U_e(1)$-degeneration with which the electromagnetic interaction is connected, we have $\frac{\partial}{\partial X_\mu}\Phi^\Sigma = Q^\Sigma \frac{\partial \chi}{\partial X_\mu}$, where Q^Σ is the electric charge of the field Ψ^Σ (given in the units of $e = \sqrt{he}/\Lambda$; in the suggested scheme, the Sommerfeld fine structure constant is $e^2/hc = 1/\Lambda^2$, where $\Lambda = \sqrt{136}$, see [15]) and $\partial\chi/\partial X_\mu$ is an a gauge (non-observable) electromagnetic potential ($\frac{\partial \chi}{\partial X_\mu} dX_\mu$ is an exact differential form). The more general phase, depending on the path in $\mathbf{A}_{3,1}$ corresponds to the real electromagnetic potential A_μ. The Lagrangian of this interaction is written as $eQ^\Sigma A_\mu j_\mu^\Sigma$, where $j_\mu^\Sigma = \overline{\Psi}^\Sigma \Gamma_\mu^\Sigma \Psi^\Sigma$ is the current of the field Ψ^Σ.

In the case of $(\mathrm{SU}(2) \otimes U(1))_i$-degeneration with which the strong interaction is connected, we have $\frac{\partial \Phi^\Sigma}{\partial X_\mu} = Y^\Sigma \frac{\partial \theta}{\partial X_\mu} + \vec{i}^\Sigma \frac{\partial \vec{\theta}}{\partial X_\mu}$, where $Y^\Sigma, \vec{i}^\Sigma$ are the

[2] It is interesting to notice that, before the transition $f \to \dot{f}$, quanta f already participated in the gravitational interaction, see [2]. Quanta $\dot{f}$ have this kind of interaction too (see earlier). Hence fundamental particles participate in the double gravitational interaction.

hypercharge and isospin of the field Ψ^{Σ}, and $\vec{\theta}, \theta$ are parameters on the group $SU(2) \otimes U(1)$.

3) We have already mentioned that interactions are switched on just after the transition $f \to \dot{f}$ when the energy-momentum realized in the transition comes in the degeneration channels. Electromagnetic and gravitational interactions are switching on when there is so much space that the evolution of fields is possible. On this stage, there are already both L- and R-fields. Hereby, as follows from the evolution equations, R-fields transform under the isotopic transformations like L-fields (it follows from this that these interactions are long-range acting and therefore quanta of the corresponding fields have zero mass).

Strong interaction is switched on inside the fiber (i.e. at supersmall distances), when there is no space at all. In such a case, evolution is impossible and therefore there are no R-fields yet. In the fiber, an R-spinor may be got as a result of involution only, i.e. by means of complex conjugation of an L-spinor φ: $\varphi \to \dot{\varepsilon}\,\bar{\varphi} = \chi'$. This kind of involution is connected with the time-reflection T (see (5)).[3] Hereby if an L-spinor φ is transformed under the law $\varphi \to u\,\varphi$, $u \in (SU(2) \otimes U(1))_i$ so an R-spinor χ' is transformed under the inverse law (in this case, we read formulas from right to left), so that we have $\chi' \to u^{-1}\chi'$. We see that the strong interaction depends on the helicity of a field: the bispinor $\psi^M = \begin{pmatrix} \varphi \\ \chi' \end{pmatrix}$ is transformed by the law $\psi^M \to e^{i\gamma_5 (\theta + \frac{1}{2}\vec{\tau}\vec{\theta})}\psi^M$, where γ_5 is the Dirac matrix. It follows from here that the parameters $\theta, \vec{\theta}$ must be pseudoscalar magnitudes. Note that ψ^M is a Majorana bispinor. The transfer to Dirac bispinors will be made in details at the consideration of a proper strong interaction.

Further we will connect these parameters with the fields of η and $\vec{\pi}$ mesons. Skeletons of these fields are represented in the fiber by the scalar form $\bar{z}z$ and vector $\bar{z}\vec{\tau}\vec{z}$. At switching off interactions, η and $\vec{\pi}$ mesons have zero masses and form the 4-isovector $(\eta, \vec{\pi})$ of the group $SL_i(2, \mathbf{C})$ which is a complex extension of the unitary group $SU_i(2)$. Hereat we have $\eta^2 - \vec{\pi}^2 = 0$, i.e. $(\eta, \vec{\pi})$ is an isotropic 4-vector like $(\bar{z}z, \bar{z}\vec{\tau}\vec{z})$. Therefore, in the fiber skeletons, η and $\vec{\pi}$ have the same norm. We put $\theta = f\eta$, where f is a normalized constant with dimension cm (not losing in generality, we can write $f = F/k$ where k is the third universal constant of the theory [1] and F is a dimensionless coupling constant). By the second quantization of fields $\eta, \vec{\pi}$, the parameters $\theta, \vec{\theta}$ obey the permulation relations, hereby $\vec{\theta}$ form the irreducible representation $D(1)$ of the group $SU_i(2)$. Hence $\vec{\theta}$ transform like $\frac{1}{2}\vec{\tau}$ or $\frac{1}{2}\bar{z}\vec{\tau}\vec{z}$ and we have to put $\vec{\theta} = \frac{1}{2} f\vec{\pi}$. (Details will be considered at the discussion of the strong interaction; it should be noted that, in the suggested scheme, the skeleton of K-meson is $\eta\,\varphi_k$ or $(\hat{\pi}\,\varphi)_k$, where φ_k are additional variables and $\hat{\pi} = \vec{\tau}\vec{\pi}$).

It follows from this reasoning that, in the general case of an arbitrary Σ, we have to write $\frac{\partial \Phi^{\Sigma}}{\partial X_\mu} = f\,\Gamma_5^{\Sigma}\left(Y^{\Sigma}\frac{\partial \eta}{\partial X_\mu} + \frac{1}{2}\vec{i}^{\Sigma}\frac{\partial \vec{\pi}}{\partial X_\mu}\right)$. We see that the strong particle interaction is described in the suggested scheme by the pseudovector coupling with derivative. The interaction Lagrangian is $f\,\overline{\Psi}^{\Sigma}\Gamma_5^{\Sigma}\Gamma_\mu^{\Sigma} \times \left(Y^{\Sigma}\frac{\partial \eta}{\partial X_\mu} + \frac{1}{2}\vec{i}^{\Sigma}\frac{\partial \vec{\pi}}{\partial X_\mu}\right)\Psi^{\Sigma}$. For the Dirac field, we have $f\,\bar{\psi}\gamma_5\gamma_\mu\left(\frac{\partial \eta}{\partial X_\mu} + \frac{1}{4}\vec{\tau}\frac{\partial \vec{\pi}}{\partial X_\mu}\right)\psi$. As is known, it is a non-renormalizable interaction. Due to the non-conserving of the pseudovector current $\bar{\psi}\gamma_5\gamma_\mu\psi$, quanta of fields η and $\vec{\pi}$ acquire the masses. Hence, strong interaction is indeed short-range acting. Forestalling (considering the proper energy diagrams of η and $\vec{\pi}$ mesons), we may just now define the mass ratio $m_\pi / m_\eta = \frac{1}{4} = 0{,}25$ [4]; its experimental rate is $\frac{134{,}97}{547{,}3} = 0{,}246$, see [11].

4) We emphasize that the strong interaction theory of fundamental particles are not a consequence of the hypothetical colour quark-gluon one: we can obtain the pseudoscalar character of η, π, K-mesons beyond the frame of this model.

We have already noticed that the strong and electromagnetic interactions appear after the transition $f \to \dot{f}$ only. At the level of quanta f, there are no these interactions (because there are no quanta $\dot{f}$ yet). Moreover, before the transition, coherent states of quanta f are electrically neutral, see [1]. Proceeding

[3] The appearance of the complex conjugation operation in quantum theory wraps itself in some mystery. This operation is always appeared at the considering of quantum jumps: we speak about the quantum transition $\psi_i \to \psi_f$ but keep in mind the transition $\psi_i \to \psi_f^*$. At the evolutional transition, complex conjugation is not used. For example, the Rubi transition in atomic physics (in time) is described by the formula $\psi(t) = \cos\Omega t\,\psi_1 + i\sin\Omega t\,\psi_2$ $(\psi_1 = \psi(0),\quad \psi_2 = \psi\left(\frac{\pi}{2\Omega}\right))$. We would say that there are two quite different kinds of motion: evolution (in space-time) and involution (in a fiber). The latter is always connected with the complex conjugation operation.

[4] In the bilocal field theory, all Feynmann diagrams converge, see Part 2 of this paper.

from the conservation law of electric charge, we may assume concerning quanta $\dot{f}$ (but it is not obligatory) that the isospinor z was neutral too before the transition, i.e. it can be written in the form of $z^0_k = z\sqrt{1+|\varepsilon|^2}\begin{pmatrix}0\\1\end{pmatrix}$ (hereby $\bar{z}^0_k z^0_k = \bar{z}_k z_k$). A charged isospinor z may be obtained from z^0 by means of the special transformation $u_0 = \frac{1}{\sqrt{1+|\varepsilon|^2}} \times \begin{pmatrix}1 & \varepsilon\\ -\bar{\varepsilon} & 1\end{pmatrix}$ belonging to the degeneration group, i.e. $z = u_0 z^0$. If, in the case of z^0, the transformation $e^{iq\chi}$ belongs to the hidden summetry group, so this transformation (i.e. $\varepsilon \to e^{i\chi}\varepsilon$) belongs in the case of z to the degeneration group. Hence we may speak about some connection between strong and electromagnetic interactions: electromagnetic interaction would not exist without strong one. Electromagnetic interaction breaks the $SU_i(2)$-symmetry.

For exactness, we call the system of waves $(\dot{f}, f)$ as space waves. We call the component f of the system as ether waves. We have seen that, at the ether level, there was only the gravitational interaction. Thus, we come to the following

T h e o r e m 1. *Phase transition 'fundamental particles → ether waves f " accomplishes the switching off all interactions (excepting the gravitational one).*

This conclusion is very important for cosmology.

Theorem 1 is a consequence of separation at the subquantum level of charges Σ from the fields produced by these charges, see [1]. But, for example, the electric charges of quanta f may perceive the external electromagnetic field (like in laboratory scattering experiments, see [16]). The theoretical description of this forced interaction was considered in [17].

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