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MULTIPARTICLE FIELDS ON THE SUBSET OF SIMULTANEITY

We propose a model describing the scattering of hadrons as bound states of their constituent quarks. We build the dynamic equations for the multiparticle fields on the subset of simultaneity, using the Lagrange method, similarly to the case of "usual" single-particle fields. We then consider the gauge fields restoring the local internal symmetry on the subset of simultaneity. Since the multiparticle fields, which describe mesons as bound states of a quark and an antiquark, are two-index tensors relative to the local gauge group, it is possible to consider a model with two different gauge fields, each one associated with its own index. Such fields would be transformed by the same laws during a local gauge transformation and satisfy the same dynamic equations, but with different boundary conditions. The dynamic equations for the multiparticle gauge fields describe such phenomena as the confinement and the asymptotic freedom of colored objects under certain boundary conditions and the spontaneous symmetry breaking under another ones. With these dynamic equations, we are able to describe the quark confinement in hadrons within a single model and their interaction during the hadron scattering through the exchange of the bound states of gluons – the glueballs.

Keywords: multiparticle fields, problem of simultaneity in relativistic quantum theory, confinement of quarks and gluons, Higgs mechanism, energy-momentum conservation law in hadron processes.

1. Introduction

Probably for the first time, the idea of multiparticle fields was proposed by H. Yukawa [1–3]. H. Yukawa called these fields "nonlocal" fields. We use another term "multiparticle fields" to show the differences between our model from the model proposed by H. Yukawa. The most essential difference between the proposed model from not only the Yukawa model, but also from models on the light cone [4, 5], quasipotential models [6–8], and models with multitime probability amplitudes [9–11] is that, in our opinion, the internal variables of such fields in different inertial reference systems cannot be related to each other, whereas these variables are connected by Lorentz transformations in the said models. We have already partially explained our viewpoint in the previous article [12]. The use of multitime probability amplitudes in [9–11, 13–15], other works of this direction, and the above-mentioned works contradicts the principles of quantum theory, because it does not consider, in our opinion, the measuring instrument influence on the state of a microsystem. In more details, we explain it in work [16], where we proposed an alternative approach to ensuring the simultaneity of quantum-mechanical measurements in different reference systems, and introduce a subset of simultaneity of the Cartesian product of several Minkowski spaces. On the other hand, the existing field theories are considered in such a way that all interaction effects are reduced only to changes in the occupation numbers of the single-particle states of free particles. This leads to the fact that, in such models, when the dynamics of processes is described, the sum of energy-momenta of these one-particle states is conserved. At the same time, the energy-momentum of hadrons, but not of constituent particles, must be conserved for the processes with hadrons. The model of multiparticle fields on the subset of simultaneity proposed in this article allows us to construct a dynamic description, which is free of the mentioned problems.

2. Scalar Product on a Subset of Simultaneity

Let us consider a meson as a two-particle system consisting of the constituent quark and antiquark. The time and coordinates of the Minkowski space of the first particle will be denoted $\left(x^0_{(1)}, x^1_{(1)}, x^2_{(1)}, x^3_{(1)}\right)$, for the second particle $\left(x^0_{(2)}, x^1_{(2)}, x^2_{(2)}, x^3_{(2)}\right)$. Here, as usual, the index 0 denotes the time coordinate of the event, and 1,2,3 are the spatial coordinates. The lower indices in parentheses identify the first and second particles. The parentheses are used to distinguish these indices from the covariant coordinates of the event. The upper indices are used to denote contravariant coordinates. The Cartesian product of Minkowski spaces for two particles is an eight-dimensional linear space. Its elements can be considered as columns

$$z^a = \begin{pmatrix} x^0_{(1)} \\ x^1_{(1)} \\ x^2_{(1)} \\ x^3_{(1)} \\ x^0_{(2)} \\ x^1_{(2)} \\ x^2_{(2)} \\ x^3_{(2)} \end{pmatrix}. \tag{1}$$

We introduce a scalar product in this eight-dimensional space by the following expression:

$$\langle z|z\rangle = \frac{1}{2}\left(g^{\mathrm{Minc}}_{ab} x^a_{(1)} x^b_{(1)} + g^{\mathrm{Minc}}_{ab} x^a_{(2)} x^b_{(2)}\right). \tag{2}$$

Here, g^{Minc}_{ab} is the Minkowski tensor. The indices a and b are repeated and summed up, and each of these indices takes the value of 0,1,2,3. Then it is convenient to use the Jacobi coordinates

$$X^a = \frac{1}{2}\left(x^a_{(1)} + x^a_{(2)}\right), \quad y^a = x^a_{(2)} - x^a_{(1)}. \tag{3}$$

In view of (3), the expression for a scalar product (2) takes the form

$$\langle z|z\rangle = g^{\mathrm{Minc}}_{ab}\left(X^a X^b + \frac{1}{4} y^a y^b\right). \tag{4}$$

A condition for the subset of simultaneity in coordinates (3) reads

$$y^0 = 0. \tag{5}$$

The coordinates of a point on a subset of simultaneity are denoted by a seven-component column

$$q^a = \begin{pmatrix} X^0 \\ X^1 \\ X^2 \\ X^3 \\ y^1 \\ y^2 \\ y^3 \end{pmatrix}. \tag{6}$$

We define the scalar product on a subset of simultaneity so that it coincides with product (4) with regard for condition (5):

$$\langle q|q\rangle = g_{ab} q^a q^b, \tag{7}$$

where the metric tensor is

$$g^{ab} = \begin{pmatrix} 1 & 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & -1 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & -1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & -4 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & -4 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & -4 \end{pmatrix}. \tag{8}$$

The multiparticle field will be described by a set of field functions $\Psi_a(q) = \Psi_a(X, \mathbf{y})$. Here, X is a set of coordinates X^0, X^1, X^2, X^3, and $\mathbf{y}$ is a set of internal variables y^1, y^2, y^3. The index a enumerates different components of the field, and its range space is determined by the representation of a transformation group, which describes the transition from field functions relative to one reference system to field functions relative to another reference system. The group of matrices acts on a subset of simultaneity as follows:

$$\hat{G} = \begin{pmatrix} \Lambda^0_0 & \Lambda^0_1 & \Lambda^0_2 & \Lambda^0_3 & 0 & 0 & 0 \\ \Lambda^1_0 & \Lambda^1_1 & \Lambda^1_2 & \Lambda^1_3 & 0 & 0 & 0 \\ \Lambda^2_0 & \Lambda^2_1 & \Lambda^2_2 & \Lambda^2_3 & 0 & 0 & 0 \\ \Lambda^3_0 & \Lambda^3_1 & \Lambda^3_2 & \Lambda^3_3 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & R^1_1 & R^1_2 & R^1_3 \\ 0 & 0 & 0 & 0 & R^2_1 & R^2_2 & R^2_3 \\ 0 & 0 & 0 & 0 & R^3_1 & R^3_2 & R^3_3 \end{pmatrix}. \tag{9}$$

The indices of the G^a_b matrix take the values from 0 to 6. $\Lambda^a_b, a, b = 0, 1, 2, 3$ are the elements of the Lorentz transformation matrix, and $R_b^a, a, b = 1, 2, 3$ are the elements of the rotation matrix.

The scalar product (7) with the metric tensor (8) is invariant relative to the group transformations (9).

Hence, our further aim will be to construct a quantum field theory not on the Minkowski space with the Lorentz group, but on the above subset of simultaneity with group (9). In work [16], we show that if the Minkowski space is replaced by a subset of simultaneity and the Lorentz group is group (9), then such a theory can be constructed in the same way as a "usual" one-particle field theory. At the same time, such a model conforms to the principle of relativity.

3. Lagrangian of a Two-Particle Meson Field

We use the notation $\psi_{c_1c_2,f_1,f_2}(q)$ for a two-particle meson field, which describes, after the quantization, the processes of creation and annihilation of bound states of a quark and an antiquark. Here, q is a set of seven variables (6). Indices with subindices 1 and 2 correspond to an antiquark and a quark, respectively, c_1 is the color of an antiquark, and c_2 is the color of a quark, f_1 is the flavor of an antiquark, and f_2 is a flavor of a quark. Accordingly, the field $\psi_{c_1c_2,f_1,f_2}(q)$ takes the value, for which the mixed tensor representations of the $SU_c(3)$ and $SU_f(3)$ groups are realized:

$$\begin{aligned}&{\psi'}_{c_1c_2,f_1,f_2}(q) = \\ &= u^{(c)\dagger}_{c_1c_3}\, u^{(c)}_{c_2c_4}\, u^{(f)\dagger}_{f_1f_3}\, u^{(f)}_{f_2f_4}\, \psi_{c_3c_4,f_3,f_4}(q).\end{aligned} \tag{10}$$

Here, $u^{(c)}_{c_2c_4}$ are the elements of an arbitrary matrix of the $SU_c(3)$ group and $u^{(f)}_{f_2f_4}$ are elements of an independent matrix of the $SU_f(3)$ group. A sign † is used to denote the elements of the adjoint matrix. Duplicate indices usually mean the summation. The dynamic equations for the field $\psi_{c_1c_2,f_1,f_2}(q)$ must be symmetric relative to transformations (10).

Moreover, the dynamic equations must be symmetric relative to group (9). The simplest Lagrangian that generates such equations can be written in the form

$$\begin{aligned}&L^{(0)} = g^{ab}\frac{\partial\psi^*_{c_1c_2,f_1,f_2}(q)}{\partial q^a}\frac{\partial\psi_{c_1c_2,f_1,f_2}(q)}{\partial q^b} - \\ &- M_\mu^2\psi^*_{c_1c_2,f_1,f_2}(q)\,\psi_{c_1c_2,f_1,f_2}(q).\end{aligned} \tag{11}$$

Here, g^{ab} are the tensor components (8), and the term M_μ will be considered as the "bare" meson mass. The "real" meson mass was considered in [16].

Since the field $\psi_{c_1c_2,f_1,f_2}(q)$ must describe the dynamics of the bound states of a quark and an antiquark, Lagrangian (11) is obviously incomplete, because it does not involve the interaction between a quark and an antiquark, which ensures the existence of a bound state. As usual, such an interaction can be introduced, if we demand the symmetry of the Lagrangian relative to the local transformations of the internal symmetry in the form (10). Since the existence of a meson as a bound state of the quark and the antiquark is due to the strong interaction, we choose the symmetry relative to the local $SU_c(3)$-transformations. This symmetry can also be achieved in the usual way, if we will replace the "ordinary" derivatives in Lagrangian (11) by the covariant derivatives and will introduce the corresponding compensating fields $A^{(1)}_{a,g_1}(q)$ and $A^{(2)}_{a,g_1}(q)$.

Further, instead of these fields, it would be convenient to consider their linear combinations, similarly to Jacobi variables,

$$\begin{aligned}&A^{(+)}_{a,g_1}(q) = \frac{1}{2}\left(A^{(1)}_{a,g_1}(q) + A^{(2)}_{a,g_1}(q)\right), \\ &A^{(-)}_{a,g_1}(q) = A^{(2)}_{a,g_1}(q) - A^{(1)}_{a,g_1}(q).\end{aligned} \tag{12}$$

A local $SU_c(3)$ group representation is given for the domain of values of the field functions $\psi_{c_1c_2,f_1,f_2}(q)$. So, this domain may be decomposed into a direct sum of subspaces which are invariant relative to transformations of this representation. Since the hadron is colorless, we will be interested in a field that has a nonzero projection only on a subspace, on which a scalar irreducible representation is realized. This means that the field $\psi_{c_1c_2,f_1,f_2}(q)$ can be given as

$$\psi_{c_1c_2,f_1,f_2}(q) = \delta_{c_1c_2}\psi_{f_1,f_2}(q), \tag{13}$$

where $\psi_{f_1,f_2}(q)$ are the new field functions for the dynamical equations, which should describe, after the quantization, the processes of creation and annihilation of mesons. These dynamic equations can be obtained from the Lagrangian with covariant derivatives that is formed, if we substitute (13) with regard for notation (12). After these transformations, this Lagrangian takes the form

$$\begin{aligned}&L_\mu = 3g^{ab}\left(\partial\psi^*_{f_1,f_2}(q)/\partial q^a\right)\left(\partial\psi_{f_1,f_2}(q)/\partial q^b\right) + \\ &+ V(q)\psi^*_{f_1,f_2}(q)\psi_{f_1,f_2}(q) - \\ &- 3M_\mu^2\psi^*_{f_1,f_2}(q)\psi_{f_1,f_2}(q),\end{aligned} \tag{14}$$

where

$$V(q) = 2g^2 g^{ab} A^{(-)}_{a,g_1}(q)\, A^{(-)}_{b,g_1}(q). \tag{15}$$

4. Dynamic Equation for the Field $V(q)$

In order to obtain the dynamic equations for a two-gluon field, we consider the simplest tensor that can be formed from single-gluon fields

$$A_{ab,g_1g_2}(q) = g^2\left(A^{(-)}_{a,g_1}(q)\, A^{(-)}_{b,g_2}(q)\right), \quad a,b = 4,5,6. \tag{16}$$

Extending the linear space of the tensors $A_{ab,g_1g_2}(q)$ relative to group (9) into the direct sum of invariant subspaces, we pick a term corresponding to the projection on a scalar subspace

$$A_{ab,g_1g_2}(q) = -A_{g_1g_2}(q)\, g_{ab} + \dots . \tag{17}$$

Convolving both sides of equality (17) with the metric tensor g^{ab}, we obtain

$$A_{g_1g_2}(q) = \frac{4}{7} g^2 \sum_{b=4}^{6}\left(A^{(-)}_{b,g_1}(q)\, A^{(-)}_{b,g_2}(q)\right). \tag{18}$$

Then we apply a similar procedure for internal indices. Considering the coupling equations obtained in [16] and definition (15), we get

$$\begin{aligned} &A_{g_1g_2}(q) = A(q)\,\delta_{g_1g_2} + \dots, \\ &A(q) = \frac{1}{14} g^2 \sum_{b=4}^{6}\left(A^{(-)}_{b,g_1}(q)\, A^{(-)}_{b,g_1}(q)\right) = \frac{1}{14} V(q). \end{aligned} \tag{19}$$

The kinetic part of the Lagrangian for the $A_{g_1g_2}(q)$ field can be given as

$$\begin{aligned} L_G^{(0)} &= \frac{1}{2} g^{ab} \frac{\partial A_{g_1g_2}(q)}{\partial q^a} \frac{\partial A_{g_1g_2}(q)}{\partial q^b} - \\ &- \frac{1}{2} M_G^2 A_{g_1g_2}(q)\, A_{g_1g_2}(q). \end{aligned} \tag{20}$$

Replacing ordinary derivatives by covariant ones and performing some calculations described in [16], we obtain the Lagrangian

$$\begin{aligned} L_V &= \frac{1}{2} g^{ab} \frac{\partial V(q)}{\partial q^a} \frac{\partial V(q)}{\partial q^b} + \\ &+ \frac{3}{2}(V(q))^3 - \frac{1}{2} M_G^2 (V(q))^2. \end{aligned} \tag{21}$$

Having a Lagrangian for the field $V(q)$, we can obtain a dynamic equation for this field such as the Euler–Lagrange equation:

$$-g^{ca}\frac{\partial^2 V(q)}{\partial q^c \partial q^a} - M_G^2 V(q) + \frac{9}{2}(V(q))^2 = 0. \tag{22}$$

We introduce the function $V(q) = V(X,\mathbf{y})$ (with regard for (6)) in the form

$$\begin{aligned} &V(X,\mathbf{y}) = V_0(\mathbf{y}) + V_1(X,\mathbf{y}), \\ &V_1(X,\mathbf{y}) \equiv V(X,\mathbf{y}) - V_0(\mathbf{y}). \end{aligned} \tag{23}$$

Then the function $V_0(\mathbf{y})$, will enter the complete Lagrangian as the potential energy of interaction of nonrelativistic constituent quarks. At the same time, it will satisfy the equation

$$4\Delta_{\mathbf{y}} V_0(\mathbf{y}) - M_G^2 V_0(\mathbf{y}) - \frac{9}{2}(V_0(\mathbf{y}))^2 = 0. \tag{24}$$

Analyzing the properties of the solutions of Eq. (24), we can obtain information about the interaction potential for quarks. Before analyzing these properties, we will make this equation to be dimensionless.

Let us introduce the dimensionless internal coordinates $\mathbf{r}$, dimensionless glueball mass m_G, and dimensionless potential energy $u(\mathbf{r})$:

$$\begin{aligned} &\mathbf{y} = l\mathbf{r}, M_G = l^{-1} m_G, \\ &V_0(\mathbf{y}) = V_0(l\mathbf{r}) = l^{-2} u(\mathbf{r}). \end{aligned} \tag{25}$$

Then, instead of Eq. (24), we obtain

$$4\Delta_{\mathbf{r}} u(\mathbf{r}) - m_G^2 u(\mathbf{r}) - \frac{9}{2}(u(\mathbf{r}))^2 = 0. \tag{26}$$

Here, $\Delta_{\mathbf{r}} \equiv \sum_{b=1}^{3} \frac{\partial^2}{\partial (r^b)^2}$ is the Laplace operator in dimensionless variables $\mathbf{r}$.

We now consider the properties of a spherically symmetric solution of Eq. (26). In order to transform the variables $\mathbf{r}\left(r^1, r^2, r^3\right)$, we pass to spherical coordinates and make the standard replacement

$$u(r) = \frac{\chi(r)}{r}. \tag{27}$$

Finally,we obtain

$$\frac{d^2\chi(r)}{dr^2} = \frac{9}{8}\frac{\chi(r)\left(\chi(r) + \left(m_G^2/9\right) r\right)}{r}. \tag{28}$$

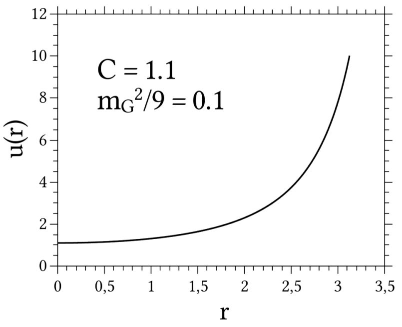

Fig. 1. Results of the numerical calculation of the dimensionless inter-quark potential $u(r)$ as a function of the dimensionless distance r for $C = 1.1$, $m_G^2/9 = 0.1$

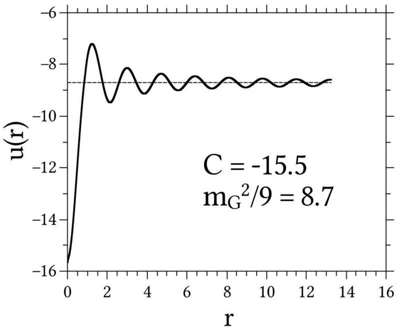

Fig. 2. Results of numerical calculations of the dimensionless inter-quark potential $u(r)$ as a function of the dimensionless distance r for $C = -15.5$, $m_G^2/9 = 8.7$

In order to analyze the properties of solutions of Eq. (28), we use an analogy with classical mechanics. We will consider the independent variable r as an analog of the time. We will call the quantity χ a "coordinate". Let its first derivative $d\chi/dr$ be a "velocity," and let the second derivative $d^2\chi/dr^2$ be an "acceleration". The dependence of "acceleration" on "coordinate", which is determined by the right part of Eq. (28), leads to the fact that, on the coordinate plane (r, χ), there are three domains [16]. Inside each of them, the "acceleration" has a constant sign. So, if the graph $\chi(r)$ gets into one of these three selected domains, then the following path of this graph is determined by the corresponding sign of the "acceleration".

Let us establish the boundary conditions for the function $\chi(r)$. We can see from Eq. (27) that if we want to obtain the finite potential energy $u(r)$ for all finite values r, we should fulfill the condition

$$\chi(r)|_{r=0} = 0. \qquad (29)$$

At that, the "initial velocity" should not be equal to zero, and we can set it to a certain real number:

$$\left.\frac{d\chi(r)}{dr}\right|_{r=0} = C, \quad C \in \mathbb{R}. \qquad (30)$$

We now consider the properties of a solution of Eq. (28) depending on the selection of the value C.

Let the solution satisfy the boundary conditions (29) and (30) with $C > 0$.

In Fig. 1, we see that, as r increases, the inter-quark potential $u(r)$ tends to infinity. Consequently, the considered model describes the quark confinement.

If $C < 0$, the potential $u(r)$ tends to some negative constant value. Thus, the eigenvalue of the squared internal Hamiltonian will definitely be negative. Since this eigenvalue is a coefficient at the squared field describing the bound state of two gauge bosons, this corresponds to the mechanism of spontaneous symmetry breaking. In this case, the result of numerical calculations of the $u(r)$ dependence on r is presented in Fig. 2.

5. Conclusions

In the proposed model, the strong interaction between the quarks in hadrons can be caused by the exchange of the bound states of gluons – the glueballs. The field $V(X, \mathbf{y})$, according to glueballs, can be represented as a sum of two terms, $V_0(\mathbf{y})$ and $V_1(X, \mathbf{y})$. The field $V_0(\mathbf{y})$ is not quantized and describes the strong interaction of quarks and gluons inside mesons and glueballs. This field satisfies the dynamic equation which describes the confiment of quarks and gluons under certain boundary conditions and spontaneous symmetry breaking – under another ones. When the bare mass of a glueball has a zero value, all solutions of this equation, irrespective of the boundary conditions, will lead to the confiment. The field $V_1(X, \mathbf{y})$ can be quantized. Though we did not consider the quantization procedure for multiparticle fields in this work, it is not different from the procedure described in work [17]. The operators obtained after the quantization will describe the processes of creation and annihilation of glueballs, as shown in [17]. Accordingly, the considered meson field quantization leads to the operators of creation and annihilation of the mesons. The meson interaction due to the interaction of constituent quarks can be described as

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the exchange by scalar glueballs. This approach differs from the one-particle field approach, because, in our model, the energy-momentum conservation law holds true precisely for hadrons, and not for the constituent particles.

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БАГАТОЧАСТИНКОВІ ПОЛЯ НА ПІДМНОЖИНІ ОДНОЧАСНОСТІ

Р е з ю м е

В роботі пропонується модель для опису процесів розсіяння гадронів як зв'язаних станів конституентних кварків. На підмножині одночасності розглядається побудова динамічних рівнянь для багаточастинкових полів за допомогою методу Лагранжа, аналогічно тому, як це робиться для "звичайних" одночастинкових полів. Розглянуто калібрувальні поля, які відновлюють локальну внутрішню симетрію на підмножині одночасності. Для багаточастинкових полів, що описують мезони як зв'язані стани кварка і антикварка і є двоіндексними тензорами відносно локальної калібрувальної групи, запропоновано модель з двома різними калібрувальними полями, кожне з яких пов'язане зі своїм індексом. Такі поля перетворюються за однаковим законом при локальному калібрувальному перетворенні і задовольняють однаковим динамічним рівнянням, але на них накладаються різні крайові умови. При певних крайових умовах ці рівняння описують такі фізичні явища, як конфайнмент і асимптотичну свободу кольорових об'єктів, а при інших крайових умовах – механізм спонтанного порушення симетрії. Ці динамічні рівняння дозволяють в межах однієї й тієї ж моделі описати як утримання кварків всередині гадронів, так і їх взаємодію в процесах розсіяння гадронів, шляхом обміну зв'язаними станами глюонів – глюболами.