Vol. 31 (2000) ACTA PHYSICA POLONICA B No 10–11

SEPARATION OF MESON AND QUARK–GLUON DEGREES OF FREEDOM IN HADRON INTERACTIONS*

A.I. Machavariani

Joint Institute for Nuclear Research, Dubna
Moscow region 141980, Russia
and High Energy Physics Institute of Tbilisi State University
University str. 9, Tbilisi 380086, Georgia

(Received July 13, 2000)

According to the conventional approach, hadrons can be treated as bound states of quarks. In this formulation one builds creation or annihilation operators of hadrons through the interacting quark operators in the Heisenberg picture $q(x)$. This approach gives another possibilities to distinguish effects generated by quark–gluon degrees of freedom.

PACS numbers: 13.75.–n, 14.70.Dj

The unique self-consistent construction of the bound states in quantum field theory was presented by Haag, Nishijima and Zimmermann in Ref. [1–3] and afterwards was developed by Huang and Weldon [4]. The nucleon annihilation operator $\mathcal{B}^{\mathrm{in(out)}}(\boldsymbol{p})$ can be constructed as a three-quark bound state annihilation operator in the following way

$$\mathcal{B}^{\mathrm{in(out)}}(\boldsymbol{p}) = \lim_{X^0\to-\infty(+\infty)} \mathcal{B}_{\boldsymbol{p}}(X^0), \tag{1}$$

where $p=(\sqrt{\boldsymbol{p}^2+m_N^2},\boldsymbol{p})$ is four momentum of the asymptotic nucleon and the Heisenberg operator $\mathcal{B}_{\boldsymbol{p}}(X^0)$ is defined through a nonlocal field operator $\Upsilon_p(X)$ of nucleon

$$\mathcal{B}_{\boldsymbol{p}}(X^0) = \int d^3\boldsymbol{X} \exp{(ipX)}\overline{u}(\boldsymbol{p})\gamma_0\Upsilon_p(X). \tag{2}$$

* Presented at the Meson 2000, Sixth International Workshop on Production, Properties and Interaction of Mesons, Cracow, Poland, May 19–23, 2000.

The composite nucleon field operator $\Upsilon_p(X)$ (which is nonlocal because it depends on the nucleon four-momentum p) is constructed through a three Heisenberg quark fields $q_i(x_i)$ with corresponding mass m_i $(i=1,2,3)$ [1–3]

$$\Upsilon_p(X) = \lim_{\rho_{12}\to 0}\lim_{\rho_3\to 0}\frac{T(q_1(x_1)q_2(x_2)q_3(x_3))}{\chi_p(X=0,\rho_{12},\rho_3)}, \tag{3a}$$

where

$$\chi_p(X,\rho_{12},\rho_3) \equiv \chi_p(x_1,x_2,x_3) = \langle 0|T(q_1(x_1)q_2(x_2)q_3(x_3))|\boldsymbol{p}\rangle \tag{4}$$

is a three quark–nucleon bound state wave function which is unambiguously determined as the solution of the bound state Bethe–Salpeter equation with a quark–gluon interaction potential, $X=(m_1x_1+m_2x_2+m_3x_3)/(m_1+m_2+m_3)$ and $\rho_{12}=x_1-x_2$, $\rho_3=(m_1x_1+m_2x_2)/(m_1+m_2)-x_3\equiv x_{12}-x_3$ denotes the center of mass and the two relative Jacobi four-coordinates.

In the other approach [4], the composite nucleon field operator is defined as

$$\Upsilon_p(X) = \int d^4\rho_{12}d^4\rho_3\chi_p^{\dagger}(X=0,\rho_{12},\rho_3)T(q_1(x_1)q_2(x_2)q_3(x_3). \tag{3b}$$

In the same manner one can also construct composite meson fields through the quark–antiquark operators.

In Refs. [1–4] it was shown that asymptotic field operators $\mathcal{B}^{\mathrm{in(out)}}(\boldsymbol{p})$ (1) satisfy the same anticommutation relations as the ordinary local field operators $b^{\mathrm{in(out)}}(\boldsymbol{p})$ for nucleons in the conventional quantum field theory. These relations make it possible to construct any "in" or "out" states with an arbitrary number of on-mass shell non-interacting, free particles. From the general standpoint, we can assume that the one particle field operators $\mathcal{B}^{\mathrm{in(out)}}(\boldsymbol{p})$ and $\mathcal{B}^{\mathrm{in(out)}}(\boldsymbol{p})$ are identical

$$\mathcal{B}^{\mathrm{in(out)}}(\boldsymbol{p}) = b^{\mathrm{in(out)}}(\boldsymbol{p})\,. \tag{5}$$

Otherwise, we must provide one hadron annihilation operator $\mathcal{B}^{\mathrm{in(out)}}(\boldsymbol{p})$ with some distinctive features which allows us to separate it from the ordinary hadron operator $b^{\mathrm{in(out)}}(\boldsymbol{p})$. This supposition contradicts with the principles of the construction of the hadronic states from the quark degrees of freedom. Therefore, we admit the validity of equation (5). Consequently, we obtain that the Fock space, the completeness condition for the asymptotic "in" or "out" real hadron fields $\sum_n |n;\mathrm{in(out)}\rangle\langle\mathrm{(out)in};n| = 1$ and even the scattering S-matrix, between an arbitrary n and m-particle states $S_{mn}=\langle\mathrm{out};m|n;\mathrm{in}\rangle$, are identical in the formulations with and without quark degrees of freedom.

Equation (5) also has one more interpretation. Thus from (5) it follows that

$$\begin{aligned}\mathcal{B}^{\rm out}(\boldsymbol{p}) - \mathcal{B}^{\rm in}(\boldsymbol{p}) &= \int d^4X {\rm e}^{ipX}\Big(i\gamma^\mu \nabla_{X^\mu} - m_N\Big)\Upsilon_p(X) \\ &= b^{\rm out}(\boldsymbol{p}) - b^{\rm in}(\boldsymbol{p}) = \int d^4x {\rm e}^{ipx}\Big(i\gamma^\mu\partial_{x^\mu} - m_N\Big)\Psi(x)\,, \end{aligned} \tag{6}$$

where $\Psi(x)$ is the local field operator of nucleon without quarks.

From relation (6) we see that the same difference between "out" and "in" operators (which defines the S-matrix), can be generated by nonlocal $\Big(i\gamma^\mu\partial_{X^\mu}-m_N\Big)\Upsilon_p(X)$ and local $\Big(i\gamma^\mu\partial_x{}^\mu-m_N\Big)\Psi(x)$ source operators. Thus, if we consider the source operator of the particle as initial-boundary (input!) conditions, then relation (6) means that by different initial-boundary conditions of our problem (with or without quarks) we can obtain the same resulting S-matrix. Therefore,the employment of quark degrees of freedom is required if by description of some scattering reaction the local source operators (*i.e.* renormalzed Lagrangians constructing from the local fields with determined quantization rules) are not enough.

Now let me consider the question: which experimental observable and theoretical relation indicates directly the quark–gluon contribution apart from the hadron degrees of freedom? The answer of this question is well known in Particle Physics. There the quark degrees of freedom were linked with the broken SU(3) or SU(6) symmetries which generate famous hadron mass splitting formulas, predict the values of the magnetic momenta of baryons, provides us with an asymptotic behavior of the S-matrix elements [5, 6] *etc.*

In my talk I will consider the consequences which follows from the Haag–Nishijima–Zimmermann picture of hadrons as quark bound states. This formulation is not restricted by a particular effective quark–gluon interaction Lagrangian and any specific dynamical mechanism can be incorporated in this scheme. Moreover, from the construction of the general field-theoretical potential follows some constrains for the quark masses and quark–meson vertex functions. For example, it was shown that if the sum of constituent quark masses is equal to the nucleon mass $m_1+m_2+m_3 = m_N$, then the pure long-distance one-pion exchange interaction between ***nucleons*** vanishes [8].

Let me turn to the above question and show that the Haag–Nishijima–Zimmermann formulation provides us with additional possibilities to distinguish effects which can be generated only by quark–gluon degrees of freedom. In particular, from the general view point, the quark degrees of freedom can be found in the hadron–hadron interaction as the result of the redefinition of the transition amplitudes

$$\langle \text{out}; m|j(0)|n; in\rangle \Leftrightarrow \langle \text{out}; m|J_p(0)|n; \text{in}\rangle, \tag{7}$$

where instead of the local hadron source operator (for instance, for pi-meson $j(x) = \left(\Box_x + m_\pi^2\right)\Phi(x)$) the quark structure generates the additional dependence on the particle four-momentum in source operator (*i.e.* for pion as quark–antiquark bound state, we have $J_p(X) = \left(\Box_X + m_\pi^2\right)\Phi_p(X)$, where $\Phi_p(X)$ is built from $q(x)$ and $\overline{q}(x)$ in similar way as nucleon nonlocal operators (3a) or (3b)). For example, for the reactions with the dominance final $\pi N - N$ vertex function with and without quark degrees of freedom must be described in the different kinematical way. Thus, if the quark degrees of freedom are not taken into account, then the πNN vertex function will be reduced to the formfactors which only depend on one variable a four momentum transfer $t = (p'_N - p_N)^2$. However, if the quark degrees of freedom are important for the considered reaction with the dominance final πN transition, then the additional dependence in πNN vertex functions over the $s = (p_\pi + p_N)^2$ and $u = (p_\pi - p'_N)^2$ variables must be observed.

Other essential difference between the field-theoretical formulations with and without quark degrees of freedom appears by construction of the effective potentials in the equations for the scattering amplitude. Based on the relativistic field-theoretical spectral decomposition method for the hadron–hadron amplitude with hadrons as quark bound states, we have derived a three-dimensional Lippmann–Schwinger-type integral equation for the corresponding amplitude. The resulting covariant equations have the following attractive features [7–9]:

I. These covariant equations can be transformed analytically and unambiguously into the four-dimensional Lorentz-invariant Bethe–Salpeter equations or other field-theoretical equations. Therefore, all the results obtained on the basis of these linearized three-dimensional equation remain valid for the other field-theoretical approaches as well.

II. The use of these relativistic equations for the investigation of scattering reactions without quarks is convenient, because the corresponding effective potential (or exchange currents) are constructed from one-variable vertex functions. In all other field-theoretical approaches the effective potentials (or exchange currents) are defined by vertices depending on two or three variables.

III. It has been demonstrated that the considered equation for hadron–hadron scattering does not change its form if nucleons and mesons are considered as bound states of quarks according to the Haag–Nishijima–Zimmermann [1–4] treatment of bound states in the quantum field theory. Moreover, we have shown that the amplitude of the suggested scattering equations with quark degrees of freedom satisfies the unitarity condition and there are no problems arising from the quark–gluon exchange poles.

The effective potential of these three-dimensional covariant equation consists of two parts: on-mass shell particle exchange terms and terms which contain all off-mass shell particle exchange diagrams. This off-mass shell particle exchange part has the form of the equal time commutator between the source operator and Heisenberg field operator which in asymptotic region transforms into particle creation or annihilation operator [7–9]. In particular, for the NN scattering this equal-time commutator term reproduces *exactly* the One Boson Exchange (OBE) Bonn model of the NN potential. If we also take into account the quark degrees of freedom, then in addition to the OBE model of the NN-potential, we also will get the quark and gluon exchange diagrams [8]. Therefore, the problem with the double account of the meson and quark degrees of freedom does not appear in the suggested formulation.

From the above consideration we can conclude that if the one-variable meson–nucleon vertex functions (*i.e.* any renormalized Lagrangian model with local fields) are enough for a description of all experimental data of the πN or NN scattering equations in low and intermediate energies, then in this energy region the quark degrees of freedom are not essential and can be omitted. Otherwise, the more detailed and complete analysis allows us to estimate the importance of the quark–gluon degrees of freedom in the πN and NN scattering reaction up to 1 GeV.

I express my deep gratitude to the organizators of the workshop MESON2000 for the possibility to give this talk.

REFERENCES

[1] R. Haag, *Phys. Rev.* **112**, 669 (1958).

[2] K. Nishijima,*Phys. Rev.* **111**, 995 (1958).

[3] W. Zimmermann, *Nuovo Cim.* **10**, 598 (1958).

[4] K. Huang, H.A. Weldon, *Phys. Rev.* **D11**, 257 (1975).

[5] V.A. Matveev, R.M.Muradyan, A.N. Tavkhelidze, *Lett. Nuovo Cim.* **7**, 719 (1973).

[6] S.J. Brodsky, G.R. Ferrar, *Phys. Rev. Lett.* **31**, 1153 (1973), *Phys. Rev.* **D11**, 1303 (1975).

[7] A.I. Machavariani, *Few-Body Syst.* **14**, 59 (1993).

[8] A.I. Machavariani, A. Buchmann, A. Faessler, G.A. Emelyanenko, *Ann. Phys.* **253**, 149 (1997).

[9] A.I. Machavariani, A. Faessler, A.J. Buchmann, *Nucl. Phys.* **A646**, 231 (1999).