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DEEP-INELASTIC PHOTOPRODUCTION OF LEPTON PAIRS AND VERIFICATION OF COLOUR SYMMETRY

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1. INTRODUCTION

About ten years ago J.D.Bjorken and E.A.Paschos [1] have proposed to use the processes with real γ-quanta to determine electric charges of the partons. Without discussing the applicability of the parton model to these processes they have stressed, however, the main difficulty of the experimental study of $\gamma N \to \gamma X$ reaction-discrimination of photons from Compton scattering and from π^0-decay. Investigations of the unified models [2,3,4] with integer charges of quarks [5,6] discovered the property of the colour suppression in reactions with deeply virtual photons which cause negligible difference between QCD and the theory with broken colour symmetry. All the difference is due to the contribution of the longitudinal polarization of gluons, produced by the Higgs mechanism. Despite the natural interpretation of the value $R = \sigma_L / \sigma_T$ for the deep inelastic $eN \to eX$ process in the unified model [7], the manifestation of the colour effects in the $e^+e^- \to \mu^+\mu^-$ process and in the muon anomalous magnetic moment is now an urgent problem. The contradiction of experimental data to the integer-charge theory was noted in paper [8]. Moreover, the mixing of gauge fields corresponding to off diagonal generators of the SU(2) and SU(3) groups of the Pati-Salam model leads to a considerable yield of the colour states in the $\nu N \to \mu X$ reaction. All these arguments cause the necessity to assume a colour-states suppressing mechanism [3]. In the integer-charge model [4] for the charged current process the colour-states yield was eliminated. Then the absence of the contradiction with the experiment was demonstrated also for the $e^+e^- \to \mu^+\mu^-$ reaction and the muon anomalous magnetic moment if the current masses of gluons $m_B \simeq 0.3 \div 0.4$ GeV [9] (m_B is the current mass of a gluon directly interacting with a lepton), and the running coupling constant of the strong interaction satisfies inequality $\alpha_s(0) \gtrsim 10^2$. Note, that the running constant α_s is definite even if the transferred momentum is very small due to the nonzero gluon mass m_B, when the energy scale parameter $\Lambda < m_B$. In the region $|q^2| \ll m_B^2$ we have no theoretical restrictions on the value of α_s, because perturbation theory cannot be applied here. The study of processes with real γ-quanta for checking the colour symmetry breaking seems to be one of the most actual now. Since the photon in the unified theory is a mixture of colour singlet and octet fields and the octet component is not suppressed by the interferences with gluon, one may hope that the colour effects will be revealed completely in photon-nucleon reactions. Such investigations were carried out /10-13/ in the framework of QCD and the parton model. The necessary condition of studying of the photon-nucleon processes is the separation of the photon-parton processes from the vector dominance contribution. In the vector dominance model the photon is connected with the parton subprocess through a superposition of the vector mesons. This leads to the necessity of introducing unknown meson structure functions /14/. However, in some kinematic region the contribution in which the photon directly takes part in the parton subprocess will dominate. This contribution can be calculated by the inhomogeneous equations of Altarelli-Parisi /15/ in the lowest order of the perturbation theory without any additional assumptions. There are the jet processes and $\mu^+\mu^-$ pair production with large k_T. Indeed, $\gamma\gamma \to \rho\rho \to$ jet + jet reaction is hardly possible, a considerable part of hadrons must go out in the direction of initial photons /16/. In connection with this the experimental discovery of two-jet processes would be an important confirmation of the simple parton model.

Unfortunately, any experimental information about deep inelastic photon-nucleon reactions is absent at the present time, except for $\gamma N \to \gamma X$ process that has been studied in Stanford /17/. The theoretical investigation of the photon photoproduction was made in /7,18/ in the framework of the integer charge theory. It was shown that the most interesting is the region of high k_T^2 of the secondary photon. The comparison with experiment shows that the integer-charge model better agrees with the data than the standard QCD. Unfortunately, the photon beam energy and k_T^2 of the secondary photons are not large enough for making an unambiguous choice between alternative theories. From the point of view of the result of /7/ the groundlessness of the assertions /19,20/ that $\eta' \to \gamma\gamma$ decay contradicts the integer charged model is especially clear, because the vector meson dominance plays the main role in this process so that the effect of colour symmetry breaking can be suppressed /21/.

The first detailed investigation of the $\gamma N \to \mu^+\mu^- X$ process for large mass of $\mu^+\mu^-$-pairs was made by R.L.Jaffe /22/ in the naive parton model. A short discussion of the low invariant mass pair production is presented in /23/. Here we consider in detail the effects of the broken colour symmetry in the process $\gamma N \to \ell^+\ell^- X$ in the theory /4/.

In Sec.2 we consider the process with a low invariant mass of the pair e^+e^-. Section 3 is devoted to the investigation of photoproduction of a large mass. Crucial characteristics - for definition of quark charges are given for both cases.

In Sec.4 we study the role of Higgs bosons in the above mentioned reaction, and show that their contribution to the cross section is small. That is why the results obtained have a model-independent character.

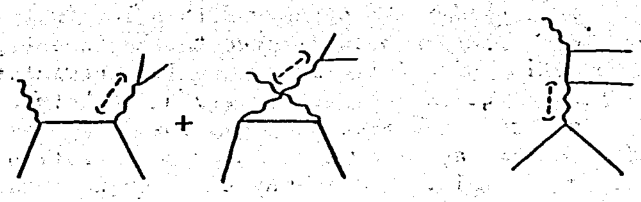
2. QUARK-GLUON SUBPROCESSES IN THE $\gamma N \to e^+e^- X$ REACTION WITH THE LIGHT LEPTON PAIR

Let us outline for completeness the main properties of the model that we have used. Unification of strong, weak, and electromagnetic interactions is realized here by mixing gauge fields corresponding to the diagonal generators of SU(2) and SU(3). All gluons become massive within the Higgs mechanism. As a consequence of unifying interactions, one of the gluons B, interacts with leptons directly. For the $q\bar{q}$ and qg reactions this leads to the effective averaging of the electric charge over the colour when the virtual momentum $|q^2| \gg m_B^2$. The quark charges are defined by matrices:

$$\begin{matrix} u \\ d \\ s \\ c \end{matrix} \begin{pmatrix} Q_1 & Q_2 & Q_3 \\ Q_1-1 & Q_2-1 & Q_3-1 \\ Q_1-1 & Q_2-1 & Q_3-1 \\ Q_1 & Q_2 & Q_3 \end{pmatrix},$$

where $Q_1 \geq Q_2 \geq Q_3$ and $\Sigma Q_i = 2$. Charges of gluons are uniquely defined by the quark charges.

Because of isotropy of the vacuum shift the masses of gluons satisfy the condition $m_B > m_{G^\pm} = m_C$ (where m_C is the mass of the second neutral gluon). For $|q^2| \ll m_W^2$ effects of weak interactions are negligible and one obtains the U(1)xSU(3) model of electromagnetic and strong interactions. The Lagrangian can be extracted from papers /9/. Let us turn now to the process $\gamma N \to \ell^+\ell^- X$ with a small invariant mass of the e^+e^- pair: $W_{e^+e^-} \ll m_B \simeq 0.3$ GeV. For this case the condition for the kinematic region, in which photons (real and virtual) are involved into parton subprocess, is the same as for the real photon photoproduction /18/. It is $k_T^2 \gg$ (1 GeV)^2. Quark subprocesses contain diagrams of two types: the Compton scattering (fig.1a) and the Bethe-Heitler process (fig.1b). In this case the co-



a) b)

Fig.1. The diagrams of parton subprocess. The broken line in brackets means the similar diagram with the B-gluon instead of the photon.

lour averaging effect due to virtual B-gluon production is absent and in both electromagnetic vertices we deal with the true electric charge of a parton. The differential cross section for the Compton type diagrams takes the form:

$$d\sigma_q^c = \frac{\alpha^3}{\pi^2}\frac{2A_Q}{s\cdot t\cdot W^2}\{\frac{\bar u}{s} + \frac{s}{\bar u} - \frac{2x^2 u_1 u_2 + (xu_1+\eta_1)(xu_2+\eta_2)}{x^2 s\bar u} + O(\frac{W^2}{s})\}(q_a(x)+\bar q_a(x))\frac{d^3p\, d^3p'}{2w\, 2w}. \qquad (2.1)$$

For the Bethe-Heitler type diagrams we get:

$$d\sigma_q^{B-H} = \frac{\alpha^3}{\pi^2}\frac{2B_Q}{st^3}\{\frac{1}{\eta_1}[x^2\bar u u_1 + xs(xu_2+\eta_2) + x(u_1-u_2)(xu_1+\eta_1) + xu_1(xu_1-xu_2+\eta_1-\eta_2)] + \frac{1}{\eta_2}[\eta_1 \to \eta_2,\ u_1 \to u_2] + O(\frac{W^2}{\eta_1}, \frac{W^2}{\eta_2})\}(q_a(x)+\bar q_a(x))\frac{d^3p}{2w}\frac{d^3p'}{2w'}. \qquad (2.2)$$

The interference terms in figs.1a and b lead to the expression:

$$d\sigma_{int} = \frac{\alpha^3}{\pi^2}\frac{C_Q}{s\cdot W^2 t^2}\{\frac{1}{s\eta_1}[t\eta_1(u_2+\frac{\eta_2}{x}) + t\eta_2(u_1+\frac{\eta_1}{x}) + xs\eta_1\bar u - (xs-\eta_1-\eta_2)(u_1(xu_2+\eta_2) + u_2(xu_1+\eta_1)) - xsu_2(xu_1+\eta_1) + 2(xu_1+\eta_1)(u_1(xu_2+\eta_2)+u_2(xu_1+\eta_1)) -$$

$$-4(s-u_1)(xu_1+\eta_1)(xu_2+\eta_2)] + \frac{1}{\bar u\eta_2}[4x^2\bar u\cdot u_1\cdot u_2 + x(u_1\eta_2-u_2\eta_1) + \bar u(\eta_1(xu_2+\eta_2)+\eta_2(xu_1+\eta_1)) - t(u_1\eta_2+u_2\eta_1) + x(xu_1+\eta_1)u_2(\bar u+s-6u_1) + (xu_2+\eta_2)(s+\bar u-2u_1)xu_1] - (\eta_1 \to \eta_2,\ u_1 \to u_2)\} \times [q_a(x)+q_a(x)]\frac{d^3p}{2w}\frac{d^3p'}{2w'}, \qquad (2.3)$$

where

$$A_Q = \overline{\sum_{1,a} Q_{1a}^4}\ ;\ B_Q = \sum_a \bar Q_a^2\ ;\ C_Q = \overline{\sum_{1,a}\bar Q_a Q_{1,a}}\ , \qquad (2.3')$$

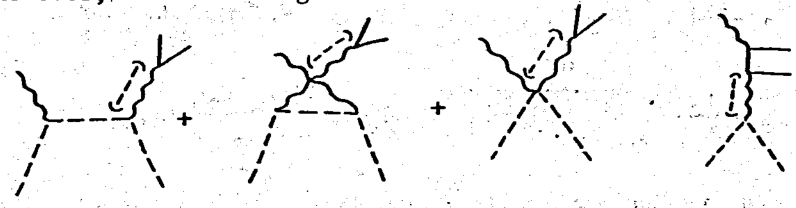
k, k' are the 4-momentum of initial and final partons, k_1 is the initial photon momentum, p, p' are the lepton momenta, ϵ, w, w' are energies of photon and leptons

$$u_{1p} = 2(pk) = xu_1;\ u_{2p} = 2(p'k) = xu_2;\ \eta_1 = 2(pk_1);\ \eta_2 = 2(p'k_1)$$

$$W^2 = 2(pp');\ s_p = 2(kk_1) = xs;\ u = \bar u - W^2;\ t = \bar t - W^2;\ \bar t = \eta_1+\eta_2;$$

$$x = \frac{t}{s-\bar u};\ \alpha = \frac{e^2}{4\pi};\ m_B \simeq m_G = \mu,$$

Direct lepton-gluon interactions that appear as a result of spontaneous symmetry breaking lead to additional gluon subprocesses represented in fig.2. The contribution of the individual diagrams there violates the unitariry of the amplitude, however, all the "dangerous" terms cancel out in the sum. The



a) b) c) d)

Fig.2

method of cancellation of the nonunitarity terms in the U-gauge was exposed in paper /7/. The Compton and B-H amplitudes have the form

$$\mathfrak{M}_C = e_\mu(k_1)\{J_\mu(k,q)\,J_\nu(q,k')\frac{1}{s-\mu^2} + J_\mu(r,k')J_\nu(k,r)\frac{1}{k'^2-\mu^2} + \epsilon_\mu(k)\epsilon_\nu(k') + \epsilon_\nu(k)\epsilon_\mu(k') - 2g_{\mu\nu}\epsilon_\beta(k)\epsilon^\beta(k')\}j_\nu, \quad (2.4)$$

$$\mathfrak{M}_{B-H} = e_\mu(k_1)J_\nu(k,k')\,\ell_{\nu\mu}(p,p'), \quad (2.5)$$

$$\ell_{\nu\mu}(p,p') = \bar{v}^+(p)\gamma_\mu\frac{-\hat{p}_1+m}{p_1^2-m^2}\gamma_\nu v^-(-p') + \bar{v}^+(p)\gamma_\nu\frac{\hat{p}_2+m}{p_2^2-m^2}\gamma_\mu v^-(-p'),$$

where

$$J_\mu(k,q) = i\{(2q_\alpha-k_\alpha)\epsilon_\alpha(k)\epsilon_\mu(q) + (2k_\alpha-q_\alpha)\epsilon_\alpha(q)\epsilon_\mu(k) - (q_\mu+k_\mu)\epsilon_\alpha(k)\epsilon^\alpha(q)\},$$

$\epsilon_\mu(k)$ is the polarization vector of gluon, $e_\mu(k_1)$ the polarization vector of photon, j_ν the lepton current. Using amplitudes (2.4), (2.5) for differential cross section we get

$$d\sigma_G^c = \frac{\alpha^3}{3\pi^2}\overline{\sum_{a,j\neq i}(Q_i-Q_j)^4}\frac{1}{s\cdot t\cdot W^2}\{6+4(\frac{s}{\bar{u}}+\frac{\bar{u}}{s})^2 - 10(\frac{\bar{u}}{s}+\frac{s}{\bar{u}}) + \frac{8}{\bar{u}s}[u_1u_2 + (u_1+\frac{\eta_1}{x})(u_2+\frac{\eta_2}{x}) + u_1(u_2+\frac{\eta_2}{x}) + u_2(u_1+\frac{\eta_1}{x})]\}(G^+(x)+G^-(x))\frac{d^3p}{2w}\frac{d^3p'}{2w}, \quad (2.6)$$

$$d\sigma_G^{B-H} = \frac{\alpha^3}{6\pi^2}\overline{\sum_{a,i\neq j}(Q_i-Q_j)^2}\frac{1}{st^3}\{\frac{1}{\eta_1}[(xs+xu-W^2)(2xu_2+\eta_2) - t\eta_2 + (2x(u_1-u_2)+\eta_1-\eta_2)(2xu_1+\eta_1)] + \frac{1}{\eta_2}[\eta_1\to\eta_2,\ u_1\to u_2]\}(G^+(x)+G^-(x))\frac{d^3p}{2w}\frac{d^3p'}{2w}, \quad (2.7)$$

$$d\sigma_{int} = \frac{2\alpha^3}{3\pi^2}\overline{\sum_{a,i\neq j}(Q_i-Q_j)^3}\frac{1}{st^2W^2}\{\frac{2(xu_2+\eta_2)}{s\eta_1}[4s\eta_1 + u_1(xs-xu_1-2\eta_1-\eta_2) - x(s-u_1)(u_1+\frac{\eta_1}{x}) - (\bar{u}+s)\eta_1] -$$

$$-\frac{1}{\eta_1}(xu_2\eta_1 + xu_1\eta_2 - 8\eta_1\eta_2) + \frac{2u_2}{\bar{u}\eta_1}[4\bar{u}\eta_1 + (xu_1+\eta_1)(s-u_1) + u_1(xs-xu_1-2\eta_1-\eta_2) - \eta_1(\bar{u}+s) + (u_1+\frac{\eta_1}{x})(xs-xu_1-2\eta_1-\eta_2)] + \frac{2u_1}{\bar{u}\eta_1}[(xu_2+\eta_2)(s-u_1 - \frac{1}{x}(2\eta_1+\eta_2))] - (\eta_1\to\eta_2)\}(G^+(x)+G^-(x))\frac{d^3p}{2w}\frac{d^3p'}{2w'}. \quad (2.8)$$

Now we should find those characteristics of the process which can discriminate the colour singlet photon from the singlet-octet mixture. For this aim we need quantities which have a little contribution both from the B-H amplitude and from the interference term. The B-H for low-mass lepton pairs is suppressed by the factor $\frac{W^2}{m_B^2} \ll 1$.

Using the symmetry of the corresponding amplitudes one can prove the absence of the interference term if one takes either $u_1 = u_2$, $\eta_1 = \eta_2$ (which are the momenta of the leptons), or integrates the cross section over lepton momenta. Quasisymmetric cross sections written in the invariant form are given in Appendix A. Transition to QCD is realized by the change $Q_i \to 2/3$ and the contribution of gluon diagrams in this case is zero. Figure 3 gives the prediction of the e^+e^- pair photoproduction on proton in the symmetrical case for the initial photon energy $\epsilon = 60$ GeV and the transversel momentum of the pair $k_T^2 = 10\ GeV^2$. The following parton distribution function is here used /24/:

$$\bar{u}(x) = \bar{d}(x) = 0{,}25\frac{(1-x)^8}{x},$$

$$G(x) = (1+9x)\frac{(1-x)^4}{x},$$

$$u(x) = \overline{u}(x) + (1-x)^3(-35{,}37T_0 - 19{,}34T_1 + 1{,}539T_2 - 0{,}7315T_3) + \frac{(1-x)^3}{\sqrt{x}}(30{,}08T_0 + 27{,}38T_1),$$

$$d(x) = \bar{d}(x) + (1-x)^4(-3{,}104T_0 - 6{,}756T_1 - 0{,}8614T_2 - $$
$$-1{,}600T_3) + \frac{(1-x)^4}{\sqrt{x}}(4{,}135T_0 + 3{,}817T_1),$$

$$T_0 = 1,\ T_1 = 2x-1,\ T_2 = 2(x\cdot 2-1)^2 - 1,\ T_3 = 4(2x-1)^3 - 3(2x-1),$$

where T_i are the Tschebyscheff polynomials.

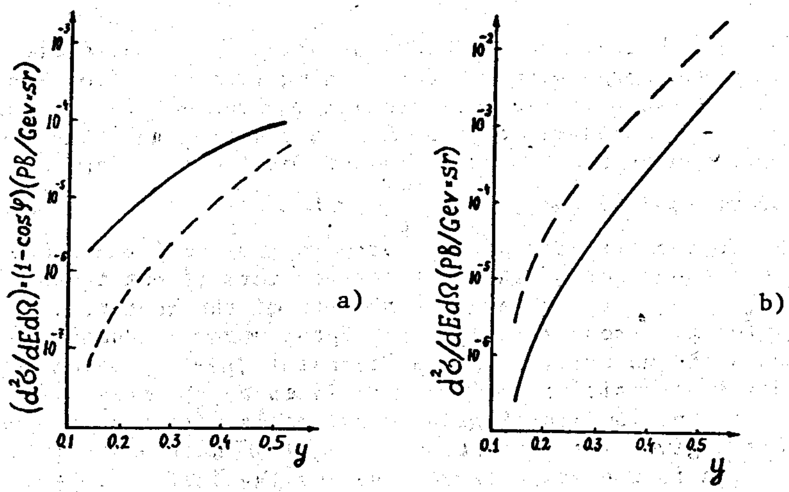


Fig.3. a) The Compton scattering diagram contribution. The solid and broken curves correspond to the quark partons and gluon partons, respectively. The ϕ is the angle between lepton momenta in the lab. frame. b) The contribution of the B-G diagrams.

Inaccuracy in fitting the distribution functions leads to significant errors in our prediction. One can considerably improve this situation using the isoscalar target and normalization to the differential cross section of $eM_{i.s.} \to eX$ reaction. Then the quark distribution functions are cancelled out similar to[18] and one obtains

$$\frac{d^2\sigma^{\gamma M_{i.s.}}}{dE\,d\Omega} \Big/ \frac{d^2\sigma^{eM_{i.s.}}}{dE\,d\Omega} = \frac{\alpha}{\pi^2 d\sigma^{eM_{i.s.}}}\{\frac{9}{5}[d\sigma_q^c + \frac{20}{9}R_d d\sigma_G^c] + d\sigma_q^{B-H} + \frac{3}{4}R_d\, d\sigma_G^{B-H}\},$$

$$\frac{d^2\sigma^{eM_{i.s.}}}{dE\,d\Omega} = \frac{5}{18}\,\frac{\alpha^2(u+\bar{u}+d+\bar{d})}{M_N\epsilon^2[1-\sqrt{1-\frac{k_T^2}{\epsilon^2y^2}}]}[1+\frac{y(1+R_d)}{(1-y)^2}(1+\sqrt{1-\frac{k_T^2}{\epsilon^2y^2}})], \tag{2.9}$$

$y = u/s$,

where the quantity

$$R_d(x) = \frac{G^+(x) + G^-(x)}{10[u(x) + \bar{u}(x) + d(x) + \bar{d}(x)]} \tag{2.10}$$

characterizes the breaking of the Callan-Gross relation and can be considered as a constant for not too small x. Experimentally $R \simeq 0.2$*. So, as is seen from fig.4, the prediction of the integer charge model for this case is more than ten times as large as that of QCD.

3. $\gamma N \to \mu^+\mu^- X$ REACTION WITH HEAVY-MASS OF LEPTON PAIR

In this case a significant contribution to $\mu^+\mu^-$ pair production comes from the B-gluon decay. Therefore, the effect of the colour symmetry breaking for photoproduction of the "heavy" lepton pair will be not so visible as compared to the previous case. Interference between amplitudes with the virtual γ and B-gluon leads to the dynamical averaging of the charge in

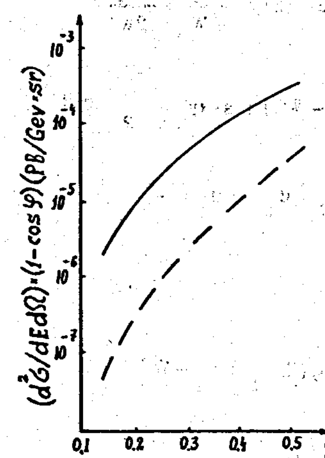


Fig.4. The differential cross section of the light-mass lepton pair production; ——— is the prediction of the integer-charge theory, – – – – the prediction of the QCD.

* As was noted[7], a nonzero value of R has natural interpretation, moreover, the functional behaviour of the R(x) distribution agrees with that of the parton. Note, that in this case the contribution of charged gluons is significant.

one of the vertices. So, expressions (2.3′) for $W^2 \gg m_B^2$ take the form:

$$A'_Q = \overline{\sum_{i,a} \bar{Q}_a^2 Q_{ia}^2}\, \frac{1}{\cos^2\phi}, \qquad C'_Q = \overline{\sum \bar{Q}_{ia}^2}/\cos^4\phi, \tag{3.1}$$

$$B'_Q = B_Q, \qquad \cos^2\phi = m_G^2/m_B^2 \simeq 1 .$$

Experimentally, however, this case is not easy for measurement, thus providing more reliable information. Besides, uncertainty due to hadronization of the photon is considerably smaller and the parton model calculations become more reliable. In the limit $W^2 \gg m_B^2$ the differential cross sections for Compton and B-H subprocesses are:

$$\begin{aligned} d\sigma^c = {} & \frac{\alpha^3}{\pi^2}\, \frac{1}{st\cdot W^2} \Big\{ 2A'_Q \Big[\frac{\bar{u}}{s} + \frac{xs}{x\bar{u} - W^2} - \frac{x^2 u_1 u_2 + (xu_1 + \eta_1)(xu_2 + \eta_2)}{xs(x\bar{u} - W^2)} + \\ & + \frac{W^2}{xs} + \frac{2W^4}{xs(x\bar{u} - W^2)} \Big] (q_a(x) + \bar{q}_a(x)) + \\ & + \frac{1}{3} \sum_{i \neq j}^{3} \overline{(Q_i - Q_j)^4}\, [xu_1(xu_2 + \eta_2 - W^2) + xu_2(xu_1 + \eta_1 - W^2) - tW^2 - \\ & - x^2 \bar{u} s + xsW^2][xs(x\bar{u} - W^2)]^{-1} (G^+(x) + G^-(x)) \Big\} \frac{d^3p}{2w}\, \frac{d^3p'}{2w'}, \end{aligned} \tag{3.2}$$

$$\begin{aligned} d\sigma^{B-H} = {} & \frac{\alpha^3}{\pi^2}\, \frac{1}{s\cdot t^3} \{ 2B_Q (q_a(x) + \bar{q}_a(x)) ([u_1(x\bar{u} - W^2) + s(xu_2 + \eta_2 - W^2) + \\ & + u_1(xu_1 - xu_2 + \eta_1 - \eta_2)] \frac{x}{\eta_1} + \frac{1}{\eta_2}[\eta_1 \to \eta_2,\ u_1 \to u_2] - \\ & - \frac{W^2 x}{\eta_1 \eta_2} [\bar{u}(x\bar{u} - W^2) + (s - \bar{u})(xs - W^2)]) + \\ & + \frac{1}{6} \sum_{i \neq j}^{3} \overline{(Q_i - Q_j)^2}\, (G^+(x) + G^-(x)) ([x(u + s)(2xu_2 + \eta_2 - W^2) - \\ & - t\eta_2 + (2xu_1 + \eta_1 - W^2)(2xu_1 - 2xu_2 + \eta_1 - \eta_2)] \frac{1}{\eta_1} + \\ & + \frac{1}{\eta_2}[\eta_1 \to \eta_2,\ u_1 \to u_2] - \frac{W^2}{\eta_1 \eta_2} [\frac{1}{2} t^2 + x^2 u\bar{u} + xs(2x\bar{u} + \eta_1 + \eta_2 - \\ & - 2W^2) - 2xu_1(xu_2 + x\bar{u} + t - W^2) - \\ & - 2xu_2(xu_1 + x\bar{u} + t - W^2)]) \frac{d^3p}{2w}\, \frac{d^3p'}{2w'}, \end{aligned} \tag{3.3}$$

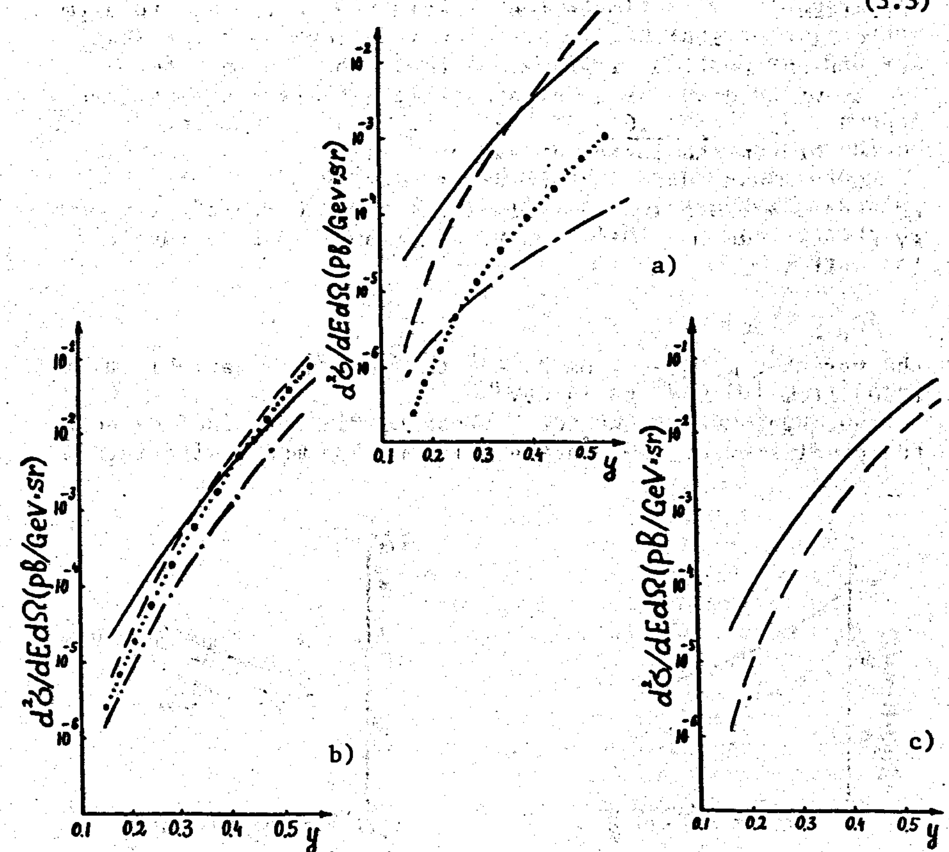


Fig.5. a) The contributions of parton subprocess for the heavy mass lepton pair ptoduction. ——— is the contribution of the Compton scattering for quark partons, – – – – the B-H process for quark partons, –·–·– and •···•··· • is the same for the gluon partons. b) The comparison of different contributions in the integer-charge theory and QCD. ——— and – – – are the Compton and B-H subprocess contributions for the integer charge, –·–·–· and •···•··· are the same for QCD. c) The predictions for the Heavy-mass lepton pair production; ——— is the integer charge theory, – – – is the QCD.

where both quark and gluon subprocesses are taken into account. Interference terms for parton subprocesses are removed just as in Sec.2. As for the B.-H. diagrams their contribution now is comparable with the Compton one which causes a further complication for the interpretation of experimental measurements. To estimate the experimental error due to asymmetrical contributions to differential cross sections (both for B-H. and Compton diagrams and their interference term), let us write down the quasi-symmetrical cross sections (see Appendix A). Figs.5,6 show the prediction for photoproduction on the proton and isoscalar targets.

Experimental data for the cross section of $\gamma N \to \mu^+\mu^- X$ process are known in the following kinematic region: the energy photon beam $\epsilon = 10 \div 11,7$ GeV, invariance mass of the lepton pair

$$4m_\mu^2 \le W^2 \le 1 \text{ GeV}^2,$$

the variable μ takes from 2 GeV^2 to 5 GeV^2, the variable u_1+u_2 takes from 104 GeV^2 to 14 GeV^2.

The magnitude of the total cross section obtained exceeds the prediction of QCD and the naive parton model with integer

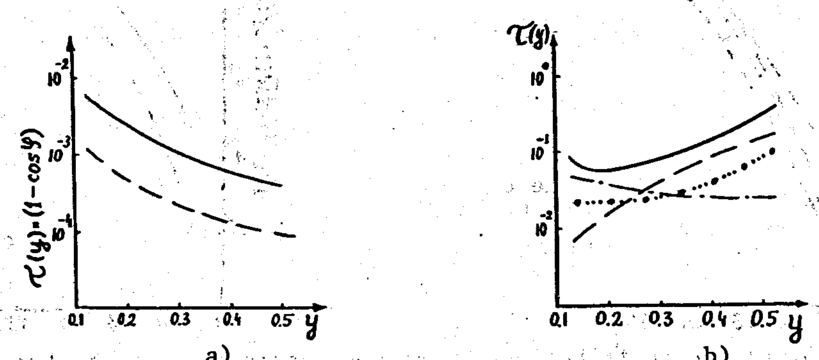


Fig.6. a) The ratio $r(y) = \dfrac{d^2\sigma^{\gamma M_{i.s.}}}{dE\,d\Omega} \Big/ \dfrac{d^2\sigma^{eM_{i.s.}}}{dE\,d\Omega}$ for an isoscalar nuclei target for light-mass lepton pair; ——— the integer-charge theory, - - - the QCD. b) The same as in Fig.6a but for heavy mass lepton pair. — — — the contribution of B-H subprocess in the integer-charge theory, —·—·— the same of Compton subprocess, ——— the sum of the both contributions, ●···●···● the prediction of the QCD.

charged quarks. One may hope that within the integer charge model the agreement of theoretical estimations with experimental data will be considerably better due to the contribution of charged gluons. Unfortunately, the measured invariant mass W is comparable with the current mass of the B-gluon. This fact prevents the application of the parton model.

4. ROLE OF HIGGS PARTICLES IN γN REACTIONS

The gauge invariance of the unified model Lagrangian necessitates the introduction of supplementary degrees of freedom associated with scalar fields. The scalar sector gives rise to some theoretical difficulties. Its source is in the zero charged behaviour of self-interaction constants of Higgs fields [2,5]. Moreover, the assumption that current masses of Higgs bosons are large [3,36] and their contribution in the experimentally accessible region is negligibly small is unacceptable in the electrostrong-interaction model. The matter is that the scalar fields are necessary for the S-matrix unitarity, i.e., in the case $m_H \gg m_B$ cross sections of all processes will grow with energy. Actually, consider, for example, the process $e^+e^- \to 3g$ (fig.7). If $m_H \gg m_B$ (fig.7c), then diagrams, fig.7a,b lead to increase of the cross section, while there is no energy for production of Higgs bosons. Experimentally, this increase of the cross section is not observed and therefore, $m_H \lesssim m_B$. Hence, we must assume that confinement (even though temporary) of scalar fields takes place, and investigations of its possible role in deep inelastic scattering are necessary. So, apriori one cannot claim that Higgs particles carry a small part of the hadron momentum and hence one can not consider partons on equal status with quarks and gluons.

Considering the scalar fields in the electrostrong model we shall use the Higgs sector that is the 3+3+3 representation of SU(3) group [9].

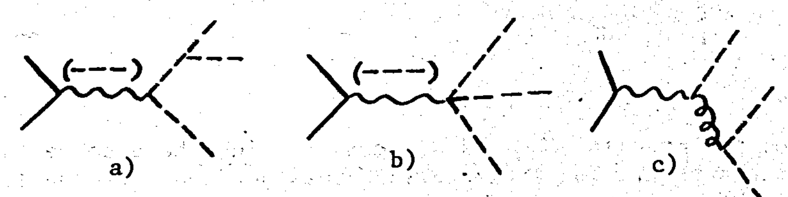


Fig.7. The Higgs boson subprocess (spiral line).

Using the experimental data

$$\int_0^1 F_2^{(p)}(x)\,dx = 0.18\pm0.01; \quad \int_0^1 F_2^{(n)}(x)\,dx = 0.12\pm0.01,$$

$$\sigma_c^{\nu}+\sigma_c^{\bar\nu} = \frac{G_F^2}{2\pi}s\cdot\frac{4}{3}(u+\bar u+d+\bar d) = \frac{G_F^2}{2\pi}s(0.475\pm0.05),$$

where F_p, F_n are the structure functions of the proton and neutron, respectively, and adding the calculation of /4/ by scalar partons, one can obtain that scalar particles transfer a small part of the nucleon momentum ($\simeq 0.02$). Therefore, for the process under investigation the contribution of scalar partons is negligible. However, they will contribute to the channels with production of scalar particles.

The part of Lagrangian is

$$\mathcal{L}^{H}_{int} = \frac{m_G e}{\sqrt6\sin\phi\cos\phi}\{\frac{2}{3}(2Q_3-Q_1-Q_2)B_\mu(\phi_D^- D_\mu^+ + \phi_D^+ D_\mu^-) +$$

$$+\frac{2}{3}(2Q_2-Q_1-Q_3)B_\mu(\phi_E^- E_\mu^+ + \phi_E^+ E_\mu^-) + \frac{2}{3}(2Q_1-Q_2-Q_3)B_\mu(\phi_V^- V_\mu^+ + \phi_V^+ V_\mu^-)\}+$$

$$+\frac{4m_G e}{3^3\sqrt6 D^2\sin\phi\cos^2\phi}\{(2Q_1-Q_2-Q_3)^3+(2Q_2-Q_1-Q_3)^3+(2Q_3-Q_1-Q_2)^3\}\phi_B B_\mu B^\mu +$$

$$+\frac{4m_G e}{9\sqrt2 D^2\sin\phi\cos^2\phi}\{(2Q_1-Q_2-Q_3)(Q_2-Q_3)+(2Q_2-Q_1-Q_3)(Q_3-Q_1)+ \tag{4.1}$$

$$+(2Q_3-Q_1-Q_2)(Q_1-Q_2)\phi_B B_\mu C^\mu + ie(A_\mu-\mathrm{ctg}\phi B_\mu)[(Q_2-Q_1)\times$$

$$\times(\phi_D^-\phi^+_{D,\mu}-\phi_D^+\phi^-_{D,\mu})+(Q_3-Q_1)(\phi_E^-\phi^+_{E,\mu}-\phi_E^+\phi^-_{E,\mu})+$$

$$+(Q_3-Q_2)(\phi_V^-\phi^+_{V,\mu}-\phi_V^+\phi^-_{V,\mu})]\},$$

$$D=\sum_i Q_i^2-\frac{1}{3}(\sum_i Q_i)^2.$$

For the process $\gamma N\to\ell^+\ell^- X$ diagrams of fig.8 have to be taken into account. The differential cross sections for the low-mass pair are

$$d\sigma^{c} = \frac{4}{81}\frac{\alpha^3}{\pi^2}\sum_a\overline{\sum_{i\neq j\neq k}(Q_i-Q_j)^2(2Q_k-Q_i-Q_j)}\frac{1}{W^2 st}(\frac{W}{m_B})^2\times$$

$$\times\{\frac{x^2u_1u_2-\eta_1\eta_2+x(u_1\eta_2+u_2\eta_1)}{x^2s^2}+\frac{u_1u_2}{\bar u^2}-$$

$$-\frac{2xu_1u_2+u_1\eta_2+u_2\eta_1}{xs\bar u}\}(G^+(x)+G^-(x))\frac{d^3p}{2w}\frac{d^3p'}{2w'}, \tag{4.2}$$

$$d\sigma^{B-H} = \frac{2\alpha^3}{\pi^2}\frac{1}{st^3}\{\frac{1}{3^4}\sum_a\overline{\sum_{i\neq j\neq k}^{3}(2Q_k-Q_i-Q_j)^2}(G^+(x)+G^-(x))+$$

$$+\frac{4}{3^8D^4}\sum_a\overline{(\sum_{i\neq j\neq k}(2Q_k-Q_i-Q_j))^2}B(x)+ \tag{4.3}$$

$$+\frac{4}{3^5D^4}\sum_a\overline{(\sum_{i\neq j\neq k}(Q_i-Q_j)(2Q_k-Q_i-Q_j))^2}C(x)\}\times$$

$$\times\{\frac{x^2}{\eta_1}[su_2+u_1(u_1-u_2)]+\frac{x^2}{\eta_2}[su_1+u_2(u_2-u_1)]\}\frac{d^3p}{2w}\frac{d^3p'}{2w'}.$$

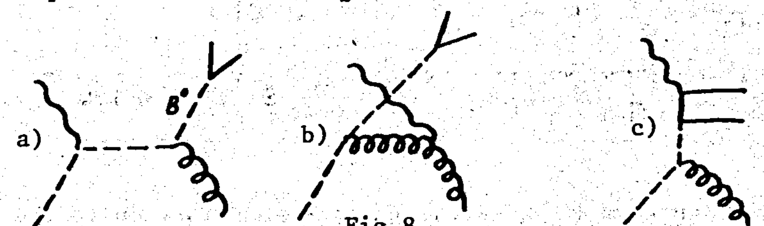


Fig.8

The contribution of the interference term for all W^2 has the form:

$$d\sigma_{int} = \frac{2\alpha^3}{3^4\pi^2}\frac{1}{st^2(W^2-m^2)}\sum_a\overline{(\sum_{i\neq j\neq k}^{3}(Q_i-Q_j)(2Q_k-Q_i-Q_j)^2)}(G^+(x)+G^-(x))\times$$

$$\times\{\frac{x}{s\eta_1}[2su_2\eta_1+u_1(2xsu_2-2xu_1u_2-u_1\eta_2-u_2\eta_1+sW^2)]+ \tag{4.4}$$

$$+\frac{xu_1}{\eta_1(x\bar u-W^2)}[xu_1(xu_1+\eta_1+x\bar u-2W^2)-\eta_1 t+xs(xu_1+\eta_1-W^2)]-\begin{pmatrix}\eta_1\to\eta_2\\ u_1\to u_2\end{pmatrix}\}\frac{d^3p}{2w}\frac{d^3p'}{2w'}.$$

One can see from (4.2)-(4.4) that parts of Compton and interference terms are suppressed by the factor $W^2/m_B^2 \ll 1$ as compared with quark and gluon subprocesses, i.e., Higgs bosons decouple from the photon.

In the case when scalar fields form a triplet, their contribution is even smaller, however, the masses of given gluon fields remain zeros in this case.

Consider now the large-mass $\mu^+\mu^-$-pair photoproduction. The cross sections in this case are

$$d\sigma^c = \frac{2a^3}{81\pi^2}\,\frac{1}{stW^2}\sum_a \overline{\sum_{i\neq j\neq k}(Q_i-Q_j)^2\,(2Q_k-Q_i-Q_j)^2}\,(G^+(x)+G^-(x)) \times$$
$$\times\Big[\frac{u_1(xu_2+\eta_2-W^2)+u_2(xu_1+\eta_1-W^2)-x\bar{u}s+W^2s}{s(x\bar{u}-W^2)}\Big]\,\frac{d^3p}{2w}\,\frac{d^3p'}{2w'}, \qquad (4.5)$$

$$d\sigma^{B-H} = \frac{2a^3}{\pi^2}\,\frac{z}{st^3}\Big\{\frac{1}{\eta_1}[su_2+u_1(u_1-u_2)]x^2+\frac{1}{\eta_2}[su_1+u_2(u_2-u_1)]x^2-$$
$$-\frac{W^2}{\eta_1\eta_2}[u_2(s-u_1)+u_1(s-u_2)]x^2\Big\}\,\frac{d^3p}{2w}\,\frac{d^3p'}{2w'}, \qquad (4.6)$$

$$z = \frac{1}{3^4}\sum_a \big(\overline{\sum_{i\neq j\neq k}(2Q_k-Q_i-Q_j)^2}\big)[G^+(x)+G^-(x)] +$$
$$+\frac{4}{3^8D^4}\sum_a\big(\overline{\sum_{i\neq j\neq k}(2Q_k-Q_i-Q_j)}\big)^2 B(x)+\frac{4}{3^5D^4}\sum_a\big(\overline{\sum_{i\neq j\neq k}(Q_i-Q_j)(2Q_k-Q_i-Q_j)}\big)^2 C(x).$$

It is clear from (4.5) (4.6) that the contribution to the cross section from Higgs bosons is insignificant because of small coefficients (~3%). A great arbitraries in the choice of scalar sector will make difficult the experimental study of this question. One needs either an idea which allows one to fix the scalar sector or a method for generating masses of vector fields without using additional scalar fields.

5. CONCLUSION

Effects caused by the colour of the photon have been clearly displayed in the photo-nucleon reactions. Theoretically, the most interesting are measurements of the yield of lepton pairs in the two cases:

a) e^+e^-- production with the invariant mass $W^2_{e^+e^-} \ll m_B^2$. The differential cross section exceeds the prediction of QCD by more than ten times. Here the role of gluons is highly essential. However, inequalities which select the region of parton subprocess imply strong requirements on the kinematical region of experiment.

b) $\mu^+\mu^-$ - production with the invariant mass $W^2 \gg m_B^2$. The differential cross section exceeds results of QCD by more than three times. The essential defect here is the large Bethe-Heitler background in the vertices of which the averaging over colour takes place.

In fact, the results have a model-independent sense because contributions from Higgs bosons in all processes are negligible. The last result is tightly connected with the fact that scalar partons carry a small part of the nucleon momentum.

In conclusion we would like to emphasize that the experimental verification of the colour symmetry seems to be one of the urgent problems of high energy physics of today, just as the discovery of the intermediate vector bosons of the electroweak interaction. The preferance of the experimental check is that the technical parameters of the photon beams existing now at SPS CERN and Fermilab seem quite suitable for this aim.

APPENDIX A

Let us write down the quasisymmetric cross section of a reaction with the light lepton pair ($\eta_2=\eta_1+\Delta$, $u_2=u_1+\delta$; $\frac{\delta}{u_1}, \frac{\Delta}{\eta_1} \ll 1$)

$$d\sigma = \frac{a^3}{\pi^2 s\cdot t\cdot W^2}\Big\{2A_Q\Big[\frac{2u_1}{s}+\frac{s}{2u_1}-\frac{x^2u_1^2+(xu_1+\eta_1)^2}{x^2u_1s}+\frac{2\delta}{s}-\frac{s\delta}{4u_1^2}+$$
$$+\frac{\delta}{2x^2su_1^2}(x^2u_1^2+(xu_1+\eta_1)^2)-\frac{\delta(2xu_1+\eta_1)+\Delta(xu_1+\eta_1)}{x^2u_1s}\Big](q_a(x)+\overline{q_a(x)})+$$
$$+\frac{1}{3}\overline{\sum(Q_i-Q_j)^4}\Big[6+4\Big(\frac{s}{2u_1}+\frac{2u_1}{s}\Big)^2-10\Big(\frac{2u_1}{s}+\frac{s}{2u_1}\Big)+ \qquad (A.1)$$
$$+\frac{7x^2u_1^2+(2xu_1+\eta_1)^2}{2x^2u_1s}+\Big(\frac{16u_1}{s^2}-\frac{10}{s}+\frac{5s}{2u_1^2}\Big)\delta-$$
$$-\frac{2\delta}{xsu_1^2}(2xu_1+\eta_1)^2+\frac{4}{xsu_1}(2\delta x(xu_1\cdot 2+\eta_1)+\Delta(2xu_1+\eta_1))\Big]\times$$
$$\times(G^+(x)+G^-(x))\Big\}\frac{d^3p}{2w}\,\frac{d^3p'}{2w'}.$$

Quantities δ and Δ are defined by experimental conditions. The contributions of the interference term is suppressed by factor $\frac{\delta}{t_1}, \frac{\Delta}{u_1}$. This is clear from expressions (2.3), (2.8) but exact formulas are too cumbersome and therefore are not written here.

The quasisymmetrical cross section of photoproduction of heavy mass muons takes the form

$$d\sigma^{c} = \frac{a^3}{\pi^2 s t W^2} \{2A'_Q[(x^2(s^2+2u^2) - \eta_1(2xu_1+\eta_1))(1-\frac{\Delta}{t}) + \frac{\delta x}{2xu_1 - W^2}(x^2(2u_1^2 - s^2) - W^2(2xu_1+\eta_1)) - \Delta(xu_1+\eta_1)] \times \frac{(q_a(x)+\overline{q_a(x)})}{xs(2xu_1 - W^2)} + \frac{1}{3}\overline{\sum_{i\neq j}(Q_i-Q_j)^4}[(2xu_1(xu_1+\eta_1-W\ ) - W^2(2\eta_1-W^2) - x^2u_1 s + xsW^2)(1-\frac{\Delta}{t}) + \Delta(u_1x - W^2) + \frac{\delta x(x^2u_1(2u_1-s)-W^4)}{2xu_1-W^2}] \times \frac{(G^+(x)+G^-(x))}{xs(2xu - W^2)}\} \frac{d^3p}{2w}\frac{d^3p'}{2w'}, \tag{A.2}$$

$$d\sigma^{B-H} = \frac{a^3}{\pi^2 s\eta_1(2\eta_1-W^2)^3}(1-\frac{\Delta}{2\eta_1 W^2})^3\{2B_Q[x(2-\frac{\Delta}{\eta})(u_1(2xu_1-W^2) + s(xu_1+\eta_1-W^2)) - \frac{W^2}{\eta_1}(1-\frac{\Delta}{\eta_1})x((xs-W^2)(s-2u_1)+2u_1(2xu_1-W^2)) + x\delta(4xu_1+xs-W^2) + \Delta xs - \frac{x^2\delta}{\eta}W^2(4u_1-s)](\overline{q_a(x)}+q_a(x)) + \frac{1}{6}\overline{\sum_{i\neq j}^{3}(Q_i-Q_j)^2}[x(2u_1+s)(2xu_1+\eta_1-W^2) - \eta_1(2\eta_1-W^2)(1-\frac{\Delta}{\eta_1}) - \frac{W^2}{\eta_1}(1-\frac{\Delta}{\eta_1})(2xs(2xu_1+\eta_1-W^2) - 2xu_1(4xu_1+4\eta_1-W^2) + (2x\delta-\Delta)(s-4u_1)x + \frac{1}{2}(2\eta_1-W^2)^2 + x\delta(xu_1-2\eta_1) + \Delta(2\eta_1-W^2))] \times (G^+(x)+G^-(x))\}\frac{d^3p}{2w}\frac{d^3p'}{2w'}. \tag{A.3}$$

Because of the small contribution from Higgs bosons, we do not write corresponding cross section in the quasi-symmetrical form.

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