CM-P00060507

Ref.TH.1919-CERN

SCALING BEHAVIOUR OF CHARGED SPIN ONE PARTONS
IN A GAUGE MODEL

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A B S T R A C T

In a hadron model in which the fermion constituents are bound by vector-isovector gauge fields, electromagnetism is introduced; by spontaneously breaking the strong (SU(2)) gauge symmetry, the gauge fields become massive. We identify the spinors and vectors with partons, and, assuming the naïve parton model hypothesis, we calculate the cross-section $e^+e^- \rightarrow \text{hadrons}$ and the structure functions of the nucleon; scaling is obtained despite the presence of an anomalous magnetic moment term in the coupling of the photon with the charged vector fields; the reason is that the spontaneous breaking of the symmetry induces a vector meson dominance-type of coupling between the photon and the neutral vector, which is just what is necessary to restore scaling.

*) Supported in part by funds of a grant of the National Science Foundation.

INTRODUCTION

Parton models ¹⁾ for hadronic structure receive a considerable attention these days. In this model, one pictures the deep inelastic electron scattering off a nucleon as if the nucleon were composed of essentially free, point-like constituents. The electron then scatters elastically and incoherently off each parton. The mechanism by which the nucleon turns into a parton state, and this into a final hadronic state is not described by the model, but it is assumed to be a consequence of some not yet understood feature of the strong interactions. For spin 0 and $\frac{1}{2}$ partons, scaling is obtained for the structure functions of the nucleon. Scaling seems to be a well-established feature of the experimental data.

In field theory the situation was unclear for some time; scaling suggests a canonical structure for the product of currents at short distances, but such a structure was not found. Recently, however, use of the renormalization group led to a demonstration of asymptotic freedom ^{2),3)} in the case of non-Abelian gauge theories. The leading high energy behaviour is determined by graphs in which the strong coupling is neglected altogether. One then obtains scaling plus logarithmic corrections. Since at present such logarithms are not of experimental interest, we will neglect them in this paper.

The identification of the parton quantum numbers is a complicated question. Feynman ¹⁾ considered the possibility of spin 0 and spin $\frac{1}{2}$ partons. We investigate in this paper the possibility of scaling behaviour of spin one partons. Standard ⁴⁾ quantum electrodynamics of massive charged spin one particles is not renormalizable, and leads to a violation of scaling ³⁾. In order to have an example of a renormalizable theory, we consider a non-Abelian gauge theory for strong and electromagnetic interactions, involving spin $\frac{1}{2}$ and spin one hadrons. This model is described in the next section. It could be extended to describe weak interactions as well, by taking $SU_L(2) \otimes SU(2) \otimes U(1)$ as a gauge group. We will not consider weak interactions, however. This can be justified by making the explicit assumption that the mass of the strong spin one parton (if they exist at all) is small as compared to the mass of the weak W 's, so that we can neglect neutral weak current effects on e^+e^- annihilation into hadrons and deep inelastic electroproduction at least at present experimental energies. In our model we have, besides the charged spin one particles, necessarily a neutral one, which couples directly to the photon, describing vector meson dominance. This coupling is completely determined by the group structure, as opposed to the "old-fashioned" ⁷⁾ vector meson dominance theory. It modifies the electromagnetic coupling to the charged vectors, by introducing effectively a "form factor".

*) This fact, as noted by Bjorken ⁵⁾ was used recently to give a phenomenological description ⁶⁾ of e^+e^- annihilation into hadrons with the help of spin one partons.

We calculate the consequences of the presence of these charged spin one partons in the nucleon, and in $e^+e^- \rightarrow \text{hadrons}$.

We conclude by considering our model in the light of existing more realistic models.

1. THE MODEL

The model is a gauge model with symmetry $SU(2) \otimes U(1)$, which was used by 't Hooft ⁸⁾ to unify electromagnetic and strong interactions. $SU(2)$ is a strong isospin gauge group, and $U(1)$ is a "hyperelectromagnetic" group. It contains an isospin $\frac{1}{2}$, hypercharge +1, doublet $N = \begin{pmatrix} p \\ n \end{pmatrix}$ of spin $\frac{1}{2}$ partons; a spin $\frac{1}{2}$ isosinglet (the electron); and a scalar, isospin $\frac{1}{2}$, hypercharge +1 doublet (the Higgs particles). The $U(1)$ gauge field is the photon Y_μ , and the $SU(2)$ gauge field is a vector-isovector $\vec{W}_\mu$ (both massless before symmetry breaking); g and g' are $SU(2)$ and $U(1)$ coupling constants.

The Lagrangian is :

$$\begin{aligned} \mathcal{L} = & -\frac{1}{4} \vec{F}_{\mu\nu} \cdot \vec{F}^{\mu\nu} - \frac{1}{4} Y_{\mu\nu} Y^{\mu\nu} + \bar{\Psi} (i \not{\partial} - m_e - g' \not{Y}) \Psi + \\ & + \bar{N} (i \not{\partial} - g \frac{\vec{\sigma}}{2} \cdot \vec{W} + \frac{g'}{2} \not{Y} - M_N) N + (D_\mu \phi)^\dagger (D^\mu \phi) - \mu^2 \phi^\dagger \phi - \\ & - h (\phi^\dagger \phi)^2 \quad (h > 0) \end{aligned} \quad (1)$$

where

$$\begin{aligned} \vec{F}_{\mu\nu} & \equiv \vec{W}_{\mu\nu} - \vec{W}_{\nu,\mu} + g \vec{W}_\mu \times \vec{W}_\nu \\ Y_{\mu\nu} & \equiv Y_{\mu\nu} - Y_{\nu,\mu} \\ N & \equiv \begin{pmatrix} p \\ n \end{pmatrix} \\ \phi & \equiv \begin{pmatrix} \phi_+ \\ \phi_0 \end{pmatrix} \\ D_\mu \phi & \equiv \left(\partial_\mu + i g \frac{\vec{\sigma}}{2} \cdot \vec{W}_\mu - i \frac{g'}{2} Y_\mu \right) \phi \end{aligned} \quad (2)$$

The commas mean derivatives; $\vec{\tau}$ are the usual Pauli matrices; Ψ is the electric field; m_e, M_N and μ are the masses of the electron, the N and the Higgs scalars, respectively.

The Lagrangian is renormalizable, and invariant under the corresponding $U(1) \otimes SU(2)$ gauge transformations.

We consider the case $\mu^2 < 0$ for which spontaneous symmetry breaking occurs as usual; the doublet ϕ develops a vacuum expectation value whose direction in isospace is spontaneously chosen; we then introduce new scalar fields with vanishing vacuum expectation value, and gauge the superfluous scalars away. Then the W fields are seen to have a common mass, whereas the renormalizability of the Lagrangian remains intact.

Defining

$$\begin{aligned} W_3^\mu &\equiv (\rho_0 - \frac{g'}{g} \gamma)^\mu \\ W_{1,2}^\mu &\equiv \rho_{1,2}^\mu \\ W_\pm^\mu &\equiv \frac{1}{\sqrt{2}} (\rho_1 \mp i\rho_2)^\mu \\ e &\equiv g' (1 + (\frac{g'}{g})^2)^{1/2} \\ g' \gamma^\mu &\equiv e A^\mu \end{aligned} \quad (3)$$

where A_μ is the photon field and e the electron's charge, then the final Lagrangian is (calling $A_{\mu\nu} \equiv F_{\mu\nu} \equiv A_{\mu,\nu} - A_{\nu,\mu}$, $\rho_{\mu\nu} \equiv \rho_{\mu,\nu} - \rho_{\nu,\mu}$ and $\rho \equiv \rho_-$, $\rho^\dagger = \rho_+$)

$$\mathcal{L} \equiv \mathcal{L}_{\text{Free}} + \mathcal{L}_{\text{Higgs}} + \mathcal{L}_{\text{VMD}} + \mathcal{L}_1 + \mathcal{L}_{\text{seagull}} + \mathcal{L}_{\text{ghost}} \quad (4)$$

where

$$\begin{aligned} \mathcal{L}_{\text{Free}} &= \frac{1}{2} \phi_{,\mu} \phi^{,\mu} - \frac{1}{2} (-2\mu^2) \phi^2 + \bar{\Psi} (i\not{\partial} - m_e) \Psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \\ &\quad - \frac{1}{2} \rho_{\mu\nu}^\dagger \rho^{\mu\nu} - \frac{1}{4} \rho_{\mu\nu}^0 \rho_0^{\mu\nu} + \frac{1}{4} g^2 \lambda^2 (\frac{1}{2} \rho_0^2 + \rho^\dagger \rho) + \bar{N} (i\not{\partial} - M_N) N \\ \mathcal{L}_{\text{VMD}} &= \frac{e}{2g} \rho_{\mu\nu}^0 F^{\mu\nu} \\ \mathcal{L}_1 &= ig [(\rho_{\mu\nu}^\dagger \rho^\mu - \rho_{\mu\nu} \rho^{\mu\dagger}) (\rho_0 - \frac{e}{g} A)^\nu + (\rho_0 - \frac{e}{g} A)_{\mu\nu} \rho^{\mu\dagger} \rho^\nu] - \\ &\quad - e \bar{\Psi} \not{A} \Psi - g\sqrt{2} (\bar{p} \not{e}_+ m + \bar{m} \not{e}_- p) + e \bar{p} \not{A} p + \\ &\quad + \frac{1}{2} g (\bar{m} \not{e}_0 m - \bar{p} \not{e}_0 p) \end{aligned}$$

This final Lagrangian is invariant under the electromagnetic gauge transformations only. It contains a lepton and hadrons; the charged fields couple to the photon; the hadrons can couple electromagnetically via the strong neutral gauge boson.

In the Weinberg model ⁹⁾, which has a gauge group similar to the one used here, one defines the ρ_0^μ (or Z^μ) and the photon A^μ to be orthogonal combinations of the gauge fields Y^μ and W_3^μ . Here we do not want to do this, to avoid a direct coupling between the electron and the strong gauge boson ρ_0^μ ; as a consequence, a direct coupling between ρ_0^μ and A^μ arises, described by $\mathcal{L}_{\text{VMD}}$. However, this is a matter of taste: if we insisted here on diagonalizing the fields, the coupling $e e \rho_0$ would be proportional to e^2 , so it is perfectly acceptable. The final results are, of course, the same.

The Lagrangian should also have a ghost part; this (as well as the Higgs and seagull parts), however, will not be of interest to us since we will not consider loops.

The Feynman rules are easily derived; they can be found in Ref. 8). We will be interested in calculations to order e^2 and zeroth order in g ; therefore the vector meson dominance graphs which have strength e/g modify, effectively, the electromagnetic current matrix elements of the $\rho^\pm$, p and n ; to lowest order, these turn out to be:

$$\begin{aligned} \langle p(p') | j_\mu(0) | p(p) \rangle &= \frac{1}{\sqrt{2E_2 E_1}} \bar{u}(\vec{p}') \gamma_\mu u(\vec{p}) \frac{\frac{1}{2} q^2 - m^2}{q^2 - m^2} \\ \langle n(p') | j_\mu(0) | n(p) \rangle &= \frac{1}{\sqrt{2E_2 E_1}} \bar{u}(\vec{p}') \gamma_\mu u(\vec{p}) \frac{\frac{1}{2} q^2}{q^2 - m^2} \\ \langle e_\pm(p') | j_\mu(0) | e_\pm(p) \rangle &= \frac{\pm 1}{\sqrt{2E_2 E_1}} \epsilon^{(\sigma')\beta*}(\vec{p}') \epsilon^{(\sigma)\alpha}(\vec{p}) V_{\mu\alpha\beta} \frac{m^2}{q^2 - m^2} \end{aligned} \quad (5)$$

where

$$\begin{aligned} q_\mu &= p'_\mu - p_\mu \\ V_{\mu\alpha\beta} &= (p+p')_\mu g_{\alpha\beta} + (q-p)_\beta g_{\mu\alpha} + (-p'-q)_\alpha g_{\mu\beta} \end{aligned} \quad (6)$$

Here, the u 's are the usual Dirac spinors and the ϵ 's are the usual polarization vectors for massive spin one particles; m is the mass of the ρ fields. We see that when $q^2 \rightarrow \infty$, the form factor of the charged spin one particles falls like $1/q^2$; this is intimately connected with the renormalizability of the theory; it is also related to the scaling behaviour in the parton model, as we will see.

2. PARTONS

We identify the p , n and ρ fields with partons in the simple parton model, rather than with the particles proton, neutron and rho.

We consider inelastic scattering of an electron by a proton, in which a final hadronic state is formed, in the one-photon exchange approximation. As is well known ¹⁾, the differential cross-section for the electron in the inclusive process is governed by the dimensionless structure functions of the nucleon $W_1(q^2, \nu)$ and $PqW_2(q^2, \nu)$; here, P and q are the four momenta of the nucleon target and the virtual photon, respectively; the variable ν is $\nu = Pq/M$, where M is the nucleon mass. The virtual total photoabsorption "cross-sections" for longitudinally or transversely polarized virtual photons are, respectively,

$$\begin{aligned}\sigma_L &= \beta \left[-W_1(q^2, \nu) + M^2 \left(1 - \frac{\nu^2}{q^2} \right) W_2(q^2, \nu) \right] \\ \sigma_T &= \beta W_1(q^2, \nu)\end{aligned}$$

where the flux factor β is conventionally chosen to be equal to $4\pi^2 \alpha (Pq + \frac{1}{2}q^2)^{-1}$.

The ratio $R \equiv \sigma_L / \sigma_T$ is then

$$R = -1 + M^2 \left(1 - \frac{\nu^2}{q^2} \right) \frac{W_2(q^2, \nu)}{W_1(q^2, \nu)} \quad (7)$$

In the simple parton model ¹⁾, the electron is assumed to scatter elastically and incoherently off essentially free, point-like constituents of the proton (partons), at high energies and momentum transfer. Assuming a sharp transverse momentum cut-off for the partons, then the four momentum p_μ of a parton is $p_\mu \simeq xP_\mu$, where x is the fraction of the longitudinal momentum of the proton that is carried by the parton, in the infinite momentum frame. The structure functions W_1 and W_2 of the proton are incoherent superpositions of the "structure functions" of the partons, $W_{1,2}^{(i)}(q^2, pq = xPq)$, which one determines from the Born graphs in electron-parton elastic scattering. These have the general form

$$W_{1,2}^{(i)}(q^2, pq) = w_{1,2}^{(s)}(q^2) \delta(q^2 + 2pq) = w_{1,2}^{(s)}(q^2) \frac{1}{2Pq} \delta(x - \frac{1}{\omega}) \quad (8)$$

which defines the functions $w_{1,2}^{(s)}(q^2)$; S is the spin of the i th parton and the variable ω is $-2Pq/q^2$.

We have then

$$\begin{aligned} W_1(q^2, \nu) &= \sum_S G_S(\omega) \left[-\frac{1}{q^2} w_1^{(S)}(q^2) \right] \\ P_q W_2(q^2, \nu) &= \frac{1}{2\omega} \sum_S G_S(\omega) w_2^{(S)}(q^2) \end{aligned} \quad (9)$$

where $G_S(\omega)$ measures the fraction of spin S partons present in the proton, with a fraction $x=1/\omega$ of the total longitudinal momentum.

We see, therefore, that scaling is realized in the parton model if and only if the limits

$$\lim_{q^2 \rightarrow -\infty} \left[-\frac{1}{q^2} w_1^{(S)}(q^2) \right], \quad \lim_{q^2 \rightarrow -\infty} w_2^{(S)}(q^2) \quad (10)$$

exist. The functions $-1/q^2 w_1^{(S)}(q^2)$ and $w_2^{(S)}(q^2)$ can be calculated, given the current matrix elements for the partons; they are positive definite and dimensionless.

We get for our model :

a) for the p partons :

$$\begin{aligned} -\frac{1}{q^2} w_1^{(1/2)}(q^2) &= \frac{1}{2} \left(\frac{\frac{1}{2} q^2 - m^2}{q^2 - m^2} \right)^2 \xrightarrow{q^2 \rightarrow -\infty} \frac{1}{8} \\ w_2^{(1/2)}(q^2) &= 2 \left(\frac{\frac{1}{2} q^2 - m^2}{q^2 - m^2} \right)^2 \xrightarrow{q^2 \rightarrow -\infty} \frac{1}{2} \end{aligned} \quad (11)$$

b) for the n partons

$$\begin{aligned} -\frac{1}{q^2} w_1^{(1/2)}(q^2) &= \frac{1}{2} \left(\frac{\frac{1}{2} q^2}{q^2 - m^2} \right)^2 \xrightarrow{q^2 \rightarrow -\infty} \frac{1}{8} \\ w_2^{(1/2)}(q^2) &= 2 \left(\frac{\frac{1}{2} q^2}{q^2 - m^2} \right)^2 \xrightarrow{q^2 \rightarrow -\infty} \frac{1}{2} \end{aligned} \quad (12)$$

c) for the $\rho_{\pm}$ partons :

$$-\frac{1}{q^2} \omega_1^{(1)}(q^2) = \frac{8}{6} \left(1 - \frac{q^2}{4m^2}\right) \left(\frac{m^2}{q^2 - m^2}\right)^2 \xrightarrow{q^2 \rightarrow -\infty} 0$$

$$\omega_2^{(1)}(q^2) = \frac{4}{6} \left(3 - \frac{q^2}{m^2} + \frac{q^4}{4m^4}\right) \left(\frac{m^2}{q^2 - m^2}\right)^2 \xrightarrow{q^2 \rightarrow -\infty} \frac{1}{6}$$
(13)

Therefore scaling is realized in the Bjorken limit.

For reference, we quote the results for standard spin 0 and $\frac{1}{2}$ partons ¹⁾

$$-\frac{1}{q^2} \omega_1^{(0)}(q^2) = 0, \quad \omega_2^{(0)}(q^2) = 2$$
(14)

and

$$-\frac{1}{q^2} \omega_1^{(1/2)}(q^2) = \frac{1}{2}, \quad \omega_2^{(1/2)}(q^2) = 2$$
(15)

which also give scaling.

It is instructive to calculate $-1/q^2 \omega_1^{(1)}(q^2)$ and $\omega_2^{(1)}(q^2)$ for the most general possible coupling between the photon and a charged spin one particle, consistent with gauge invariance, parity and time reversal symmetries, and structurelessness of the particle; here ¹⁰⁾ the vertex is :

$$V_{\mu\alpha\beta} = (p+p')_{\mu} g_{\alpha\beta} + 2\mu (q_{\beta} g_{\alpha\mu} - q_{\alpha} g_{\beta\mu}) + e (p+p')_{\mu} q_{\alpha} q_{\beta}$$
(16)

where the constants μ and e are real but arbitrary, and correspond to an intrinsic arbitrary magnetic dipole moment and electric quadrupole moment, respectively. One obtains

$$-\frac{1}{q^2} \omega_1^{(1)}(q^2) = \frac{8}{6} \mu^2 \left(1 - \frac{q^2}{4m^2}\right)$$

$$\omega_2^{(1)}(q^2) = \frac{4}{6} \left[3 - \frac{q^2}{m^2} + \frac{q^4}{4m^4} - 2\mu(\mu-1) \frac{q^2}{m^2} \left(1 - \frac{q^2}{2m^2}\right) + \right.$$

$$\left. + 2e q^2 \left(1 - \frac{3q^2}{4m^2} + \frac{q^4}{8m^4}\right) + 2\mu e \frac{q^4}{m^2} \left(1 - \frac{q^2}{4m^2}\right) + e^2 q^4 \left(1 - \frac{q^2}{4m^2}\right)^2 \right]$$
(17)

The Yang-Mills coupling corresponds to $\mu=1$, $\rho=0$, and the minimal substitution coupling (without anomalous magnetic moment ⁴⁾), to $\mu=\frac{1}{2}$, $\rho=0$. For no values of μ , ρ , scaling is obtained; this is a consequence of the "bad" high energy behaviour of the (not renormalizable) theory.

Going back to our model, we obtain asymptotically

$$\begin{aligned} W_1(q^2, \nu) &= \frac{1}{4} G_{1/2}(\omega) + O\left(\frac{m^2}{q^2}\right) \\ Pq W_2(q^2, \nu) &= \frac{1}{2\omega} \left(G_{1/2}(\omega) + \frac{1}{6} G_1(\omega) \right) + O\left(\frac{m^2}{q^2}\right) \end{aligned} \quad (18)$$

As for the ratio R , one obtains

$$R = \frac{1}{6} \frac{G_1(\omega)}{G_{1/2}(\omega)} + O\left(\frac{m^2}{q^2}, \frac{m^2}{q^2}\right) \quad (19)$$

in the Bjorken limit. A similar effect ¹⁾ (non-vanishing asymptotic R) is found if one mixes spin $\frac{1}{2}$ and 0 partons.

As opposed to the simple ¹⁾ spin 0 and $\frac{1}{2}$ parton models, we have here an explicit q^2 dependence of W_1 and PqW_2 (whose limit $q^2 \rightarrow \infty$ exists); one might therefore feel tempted to compare these functions with experimental data, at fixed ω as functions of q^2 . This, however, should not be taken seriously, because non-leading (i.e., non-parton-like behaviour) terms presumably compete with the non-asymptotic terms in W_1 and PqW_2 . In fact, if one insists in comparing with experiment, the disagreement is even qualitative, unless the spin one partons contribute appreciably at large ω only ⁶⁾.

In our model, the neutral (n) and charged (p) spin $\frac{1}{2}$ partons contribute asymptotically equal amounts to W_1 and PqW_2 , but with a strength equal to 1/4 times the standard spin $\frac{1}{2}$ charged partons. This is a consequence of their isospin contents.

We can use our model to calculate the cross-section for $e^+e^- \rightarrow \text{hadrons}$. We obtain

$$\sigma(e^+e^- \rightarrow p\bar{p}) = \frac{\pi\alpha^2}{3s} \theta(s-4M_N^2) \sqrt{1-\frac{4M_N^2}{s}} \left(1 + \frac{2M_N^2}{s}\right) \left(\frac{s-2M_N^2}{s-M_N^2}\right)^2 \quad (20.a)$$

$$\sigma(e^+e^- \rightarrow m\bar{m}) = \frac{\pi\alpha^2}{3s} \theta(s-4M_N^2) \sqrt{1 - \frac{4M_N^2}{s}} \left(1 + \frac{2M_N^2}{s}\right) \left(\frac{s}{s-M_N^2}\right)^2 \quad (20.b)$$

$$\sigma(e^+e^- \rightarrow e^+e^-) = \frac{\pi\alpha^2}{12s} \theta(s-4m^2) \left(1 - \frac{4m^2}{s}\right)^{3/2} \left(\frac{m^2}{s-m^2}\right)^2 \left(\frac{s^2}{m^4} + 20\frac{s}{m^2} + 12\right) \quad (20.c)$$

to be compared with

$$\sigma(e^+e^- \rightarrow \mu^+\mu^-) \rightarrow \frac{4\pi\alpha^2}{3s} \quad (21)$$

Asymptotically, then

$$\begin{aligned} \sigma(e^+e^- \rightarrow p\bar{p}) &= \sigma(e^+e^- \rightarrow m\bar{m}) = \frac{\pi\alpha^2}{3s} \\ \sigma(e^+e^- \rightarrow e^+e^-) &= \frac{\pi\alpha^2}{12s} \end{aligned} \quad (22)$$

so the ratio $\tilde{R}(s) = \sigma(e^+e^- \rightarrow \text{hadrons}) / \sigma(e^+e^- \rightarrow \mu^+\mu^-)$ is, asymptotically

$$\tilde{R}(\infty) = \frac{9}{16} \quad (23)$$

Our result (20.c) for spin one partons is plotted in the Figure. For the sake of plotting we put $m=1$ GeV. For comparison, we also give the predictions of some other couplings.

In the Weinberg⁹⁾ model the presence of neutrinos and thus the parity violation enters in the calculations^{*)}. Hence we do not consider the Weinberg model as a model for spin one partons. Nevertheless, it is interesting to note how model dependent the predictions for $\tilde{R}(s)$ are.

3. DISCUSSION

We think the most important conclusion to be drawn from the foregoing is a qualitative one : in our renormalizable model charged spin one partons scale. This has to be compared with the naïve minimal coupling theory. There one finds

*) Details of these calculations will be presented elsewhere.

$\tilde{R}(s) = s/4m^2$. Asymptotically this would violate unitarity¹⁰⁾; also hadronic corrections to quantum electrodynamics would become uncontrollable⁵⁾. Taking into account that in gauge theories spin one particles have an anomalous magnetic moment makes the situation worse. In fact, one finds, for a naïvely massive Yang-Mills coupling, $\tilde{R}(s) = s^2/16m^4$. The vector dominated graph then cancels two powers of s and restores scaling. Further, it is remarkable that though the neutral spin one particle gives effectively a form factor to the spin $\frac{1}{2}$ partons, there results no fall-off in their contribution to the asymptotic behaviour of the current matrix element.

It should be stressed that Eqs. (18) and (22) are supposed to hold asymptotically to all orders in g , with only logarithmic corrections. This can be shown^{2),11)} using the Callan-Symanzik equation and the fact that gauge theories become asymptotically free^{*)}. The coefficients of the logarithmic terms as well as Eq. (23) are model dependent (the latter through the inhomogeneous term in the appropriate Callan-Symanzik equation).

We have considered an extremely simple model. Many complicated models exist, none seems to be plausible¹³⁾, however. On the one hand, we have quark schemes [we should extend our gauge group to $SU(3) \otimes U(1)$ or $SU(4) \otimes U(1)$], but this would not modify the qualitative aspects of our results. Perhaps quark quantum numbers are slightly favoured (cf., the predictions for σ^{vD}/σ^{eD})¹⁴⁾ over integral charge values. In any case the spin one gauge gluons are usually coupled to colour only¹⁵⁾. This means that they have no charge and do not couple to the photon. On the other hand, theories¹⁶⁾ with charged spin one particles have been constructed and at present there seems no compelling reason to exclude these. An interesting example is the $SU(3) \otimes SU(3) \otimes U(1)$ model of Han and Nambu¹⁷⁾ with integrally charged coloured quarks.

The present investigation does not explain the failure¹⁸⁾ of scaling in electron-positron annihilation at CEA/SPEAR energies. Yet this process is interesting in the sense that all charged fundamental particles will contribute to $\tilde{R}(s)$. This has to be contrasted with deep inelastic scattering. There it is the parton wave function of the proton that is measured, rather than the existence of various partons. The factor $1/6$ in (18) and (19) will give a further suppression of spin one effects.

At any rate, considering processes in which the parton model seems to work, and abstracting the results of our Lagrangian model, there is no evidence against the existence of charged spin one partons.

*) We ignore the fact that the Higgs particles ϕ may upset the asymptotic freedom. For a discussion of this point, especially of the possibility of a dynamical symmetry breakdown, see Refs. 3),12).

The authors would like to acknowledge fruitful discussions with members of the CERN Theoretical Study Division, especially with N. Cabibbo, J. Ellis and C.H. Llewellyn Smith. One of us (M.A.F.) is further indebted to J. Primack for encouragement and interesting discussions. He would also like to thank the CERN Theoretical Study Division for hospitality while part of this work was done. The second author (G.J.K.) thanks J. Cleymans, G. 't Hooft and K. Rothe for helpful comments.

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FIGURE CAPTION

Charged spin one pair production contribution to $\tilde{R}(s)$
with different couplings :

- I) QED plus anomalous magnetic moment.
- II) "Naïve" QED coupling.
- III) The Weinberg model.
- IV) The present model.

The vector boson mass has been set to $m = 1$ GeV.