Phys. C 1, 25 (1979).

[4]See, e.g., A. Pais, in *Orbis Scientiae, High Energy Physics in the Einstein Centennial Year*, edited by A. Perlmutter and L. F. Scott (Plenum, New York, 1979), Vol. XVI.

[5]See, e.g., Ch. Berger *et al.* (PLUTO Collaboration), Phys. Lett. 81B, 410 (1979).

[6]Partial results for $\sqrt{s}$ = 27, 30, and 31.6 GeV were given previously by Berger *et al.*, Ref. 1a.

[7]Ch. Berger *et al.* (PLUTO Collaboration), Phys. Lett. 91B, 148 (1979).

[8]$\gamma\gamma \to \mu^+\mu^-$ does not survive the criteria for hadron-event selection.

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[11]R. D. Field and R. P. Feynman, Nucl. Phys. B136, 1 (1978).

[12]See A. Ali, Ref. 3.

[13]The charmed mesons have an average branching ratio of about 10% for semileptonic decays. See R. L. Kelley *et al.* (Particle Data Group), Rev. Mod. Phys. 52, No. 2, Pt. 2, S1 (1980). See also W. Bacino *et al.*, Phys. Rev. Lett. 45, 329 (1980). The assumption of 10% branching ratio for muonic decays of b and t mesons is model dependent. The prediction of muon rate varies in proportion to the assigned value of the branching ratio.

[14]See, e.g., H. Georgi and M. Machacek, Phys. Rev. Lett. 43, 1639 (1979), and references therein.

[15]The function $f(z) = 1 - z$ also falls very close to the lower edge of the bands.

Is the $x \to 1$ Behavior of the Deep-Inelastic Structure Functions in the Bjorken Limit Sensitive to the Confinement Mechanism?

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This Letter incorporates confinement to the covariant parton model and investigates its consequences on the qualitative behavior of the deep-inelastic structure functions in the Bjorken limit. It is proved that this incorporation does not change the known behavior of the baryonic structure functions, but it leads, for mesons, to a nonvanishing value of $\nu W_2(\omega^{-1})$ and $W_1(\omega^{-1})$ in the $\omega^{-1} \to 1$ limit and to a violation of the Callan-Gross relation.

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The deep-inelastic structure functions, νW_2 and W_1, associated to deep-inelastic electroproduction processes are directly related to the absorptive part of the virtual Compton amplitude $\gamma^* h \to \gamma^* h$. In the covariant parton model[1] and in the Bjorken limit, it can be seen that this amplitude is dominated by the hand-bag diagrams (Fig. 1), where the internal lines are associated with quarks.

In this model, if one applies the usual ideas of analyticity to the quark-hadron amplitude, the Bjorken scaling for $\nu W_2(\omega^{-1})$ and $W_1(\omega^{-1})$ is obtained[2] $[\omega = 2\nu/(-q^2)]$.

Azcoiti, Alonso, and Cruz[3] have proved that if, in the context of the covariant parton model, one takes into account the differences which the quark model establishes between baryons and mesons, one can obtain, in turn, essential differences between baryonic and mesonic structure functions. In particular, they proved that if the sum of the free masses of the valence quarks in a meson is less than the meson mass, $m_q + m_{\bar{q}} < M_m$, one can obtain a nonvanishing value of the mesonic structure functions in the $\omega^{-1} \to 1$ limit, as well as a violation of the Callan-Gross relation for mesons.

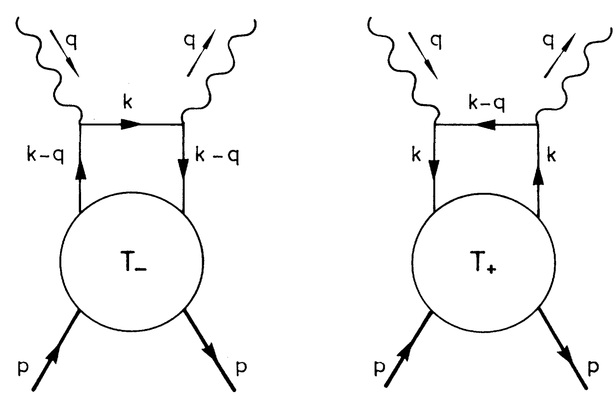

FIG. 1. Hand-bag–diagram contributions to the virtual Compton amplitude.

This special behavior of the mesonic structure functions in the $\omega^{-1}\to 1$ limit, which had been pointed out by other authors[4] based on angular momentum conservation arguments, is due to the contribution of quark single-particle intermediate states to the quark-meson amplitude. Evidently, this contribution does not appear when the hadron is a baryon.

In our earlier work,[3] we proceeded as follows: The amplitude $T_-(\mu^2,s')$ $[\mu^2=(k-q)^2, s'=(p-k+q)^2, u'=(-p-k+q)^2]$ (see Fig. 1) has singularities in the μ^2, s', and u' variables. If one denotes by $\rho(\mu^2)/(s'-m_\alpha{}^2)$ the contribution to $T_-(\mu^2,s')$ of the one-quark single-particle intermediate state in the s' channel, it is implicit in Ref. 3 the supposition that $\rho(\mu^2)$ has no cuts in the μ^2 variable. This supposition, besides a "partial" incorporation of confinement in the sense that a meson cannot decay into two quarks on mass shell, led to the result quoted above on the qualitative behavior of the mesonic structure functions.

We shall prove next, that, if we incorporate confinement to the covariant parton model in a more general way, we automatically establish that $\rho(\mu^2)$ has no singularities in the μ^2 variable and, consequently, we reproduce the results of our earlier work,[3] independently of the quark valence masses in a hadron.

Unitarity relates the total cross section $\gamma^*h\to X$ with the imaginary part of the forward $\gamma^*h\to\gamma^*h$ amplitude. At this level, we write the unitarity relation only with hadrons or, equivalently, find that hadrons are the only possible final states of the γ^*h process.

However, when we consider the quark-hadron $\to$ quark-hadron amplitude, it is necessary to incorporate quarks in the unitarity relation for this process; that is to say, *quarks are possible final states of the quark-hadron process.*

If we consider an amplitude involving only quarks and at least two hadrons, that amplitude is related, by crossing and *CPT* invariance, to the amplitude $h_1h_2\to$ quarks + hadrons. It is well known that, because of confinement, these processes do not occur in nature. Therefore, the simplest way to formulate confinement is to impose the condition that any amplitude involving only quarks and hadrons[6] must be zero *when all the external particles are on the mass shell.*

The incorporation of this fact in the covariant parton model leads to $\rho(\mu^2)$ having no singularities in the μ^2 variable. In fact, $\rho(\mu^2)$ is essentially the product of two amplitudes of the type quark-hadron $\to$ quark, where the quark in the initial state has a mass equal to μ^2 *and the quark in the final state is on the mass shell* ($s'=m_\alpha{}^2$). Now unitarity and analyticity tell us that all the singularities of $\rho(\mu^2)$ in the μ^2 variable are due to possible *on-mass-shell* intermediate states in the μ^2 channel of each of the two mentioned amplitudes defining $\rho(\mu^2)$. This implies that we are only concerned with the on-mass-shell values of amplitudes involving only quarks and hadrons, which, as a consequence of our confinement formulation, are zero; so $\rho(\mu^2)$ has no singularities.

Let us proceed under these conditions to calculate the contribution to $\nu W_2(\omega^{-1})$ of an intermediate single-particle state in the s' channel of $T_-(\mu^2,s')$.[7] We use the Sudakov variables

$$K=xp+yq+\kappa$$

with κ orthogonal to p and q, spacelike, and two dimensional. In the Bjorken limit, the dominant contribution of Fig. 1 to the structure functions comes from those regions in which μ^2 is kept finite when $q^2\to-\infty$ (as a consequence of the softness hypothesis of the covariant parton model). The only one of these regions giving nonzero contribution to $\nu W_2(\omega^{-1})$ is $y\sim 1$. If we make the change of variables $y=1+\bar{y}/2\nu$, the contribution to $\nu W_2(\omega^{-1})$ from an intermediate single-particle state in the s' channel of $T_-(\mu^2,s')$ can be written[3] as

$$\lim_{\mathrm{Bj}}\nu W_2{}^\alpha=\mathrm{Im}\left\{\frac{iM^2}{4(2\pi)^4}\int_{-\infty}^{+\infty}dx\int_{-\infty}^{+\infty}d\bar{y}\int d^2\kappa\,\frac{4x^2\rho(\mu^2)}{(x-\omega^{-1}+i\epsilon)(s'-m_\alpha{}^2+i\epsilon)}\right\}, \tag{1}$$

where

$$\mu^2=x\bar{y}+x^2M^2+\kappa^2,$$

$$s'=(x-1)\bar{y}+(x-1)^2M^2+\kappa^2,$$

$$u'=(x+1)\bar{y}+(x+1)^2M^2+\kappa^2.$$

As, in our model, $\rho(\mu^2)$ is analytic in all the complex $\bar{y}$ plane, $\rho(\mu^2)$ is either a constant or not bound. The softness hypothesis of the covariant parton model implies that $\rho(\mu^2)$ goes to zero as μ^2 becomes large and, consequently, $\rho(\mu^2)$ is not bound. Therefore, the integrand of (1) does not have the right asymptotic behavior for the vanishing of the integral over the infinite circle.

The calculation of (1) gives, as a result, two terms.[3] The first can be written as

$$-\frac{M^2\omega^{-2}}{16\pi}\int_{-\infty}^{\mu_0^2(\omega,m_\alpha,M)}\rho(\mu^2)\,d\mu^2, \tag{2}$$

where $\mu_0^2(\omega,m_\alpha,M)=\omega^{-1}M^2+m_\alpha^2\omega^{-1}-(\omega^{-1}-1)$. Because of the softness hypothesis this integral is convergent, positive $[\rho(\mu^2)<0]$, and goes to zero when ω^{-1} goes to 1.

The second contribution can be written as[3]

$$\frac{M^2}{8(2\pi)^3}\int_0^\infty d(-\kappa^2)\mathrm{P}\left[\int_{-\infty}^{+\infty}\frac{4x^2\,dx}{x-\omega^{-1}}\left(\mathrm{P}\int_{-\infty}^{+\infty}\frac{\rho(\mu^2)\,d\bar{y}}{(x-1)\bar{y}-m_\alpha^2+(x-1)^2M^2+\kappa^2}\right)\right]. \tag{3}$$

This integral is convergent as a consequence of the softness hypothesis, too, and it is responsible for the invariance violation of the Callan-Cross relation in the case of spin-$\frac{1}{2}$ quarks. This contribution takes into account the final-state-quark interaction.[3]

Similarly, one can compute the contribution νW_2^{β} to $\nu W_2(\omega^{-1})$, from an intermediate single-particle quark state in the u' channel of the quark-meson→quark-meson amplitude. If we write this contribution to $T_-(\mu^2,s')$ as

$$\frac{\eta(\mu^2)}{u'-m_\beta^2+i\epsilon}$$

after straightforward calculations,[3] we obtain two new contributions to $\nu W_2(\omega^{-1})$. The first one is similar to (2) and can be written as

$$-\frac{M^2\omega^{-2}}{16\pi}\int_{-\infty}^{\mu_0^2(-\omega,im_\beta,M)}\eta(\mu^2)\,d\mu^2, \tag{4}$$

where $\mu_0^2(-\omega,im_\beta,M)=-\omega^{-1}M^2+m_\beta^2\omega^{-1}/(\omega^{-1}+1)$. This contribution is convergent, positive $[\eta(\mu^2)<0]$, and very interesting because of its nonvanishing value in the $\omega^{-1}\to 1$ limit. If we take this limit at (4), we obtain the following *constant at* $\omega^{-1}=1$:

$$-\frac{M^2}{16\pi}\int_{-\infty}^{(m_\beta^2/2)-M^2}\eta(\mu^2)\,d\mu^2. \tag{5}$$

The second contribution to $\nu W_2(\omega^{-1})$, from the intermediate single-particle state in the u' channel of $T_-(\mu^2,s')$, is similar to (3), and can be written as

$$\frac{M^2}{8(2\pi)^3}\int_0^\infty d(-\kappa^2)\mathrm{P}\left[\int_{-\infty}^{+\infty}\frac{4x^2\,dx}{x-\omega^{-1}}\left(\mathrm{P}\int_{-\infty}^{+\infty}\frac{\eta(\mu^2)\,d\bar{y}}{(x+1)\bar{y}-m_\beta^2+(x+1)^2M^2+\kappa^2}\right)\right], \tag{6}$$

and it is also responsible for a violation of the Callan-Gross relation in the spin-$\frac{1}{2}$ quark case.

Our model adds nothing new to the calculation of the contribution to νW_2 from the intermediate multiparticle states of $T_-(\mu^2,s')$. This calculation has been made in Ref. 1, and can be written as

$$-\frac{M^2\omega^{-2}}{(2\pi)^3}\int_{\bar{y}_1}^{-\infty}d\bar{y}\int d^2\kappa\,\mathrm{Im}\overline{T}_-(\mu^2,s'), \tag{7}$$

where

$$s'=(\omega^{-1}-1)\bar{y}+(\omega^{-1}-1)^2M^2+\kappa^2,$$

$$\mu^2=\omega^{-1}\bar{y}+\omega^{-2}M^2+\kappa^2,$$

$$\bar{y}_1=\frac{s_0'}{\omega^{-1}-1}-(\omega^{-1}-1)M^2-\frac{\kappa^2}{\omega^{-1}-1},$$

and where the integral (7) extends along the cut of the s' channel of the $T_-(\mu^2,s')$ amplitude (s_0' is the beginning of this cut). $\overline{T}_-(\mu^2,s')$ is what is left after subtracting from $T_-(\mu^2,s')$ the contributions of all the possible intermediate single-particle states in the s' and u' channels. The contribution (7) to $\nu W_2(\omega^{-1})$ is convergent and goes to zero when $\omega^{-1}\to 1$, since $\bar{y}_1$ goes to $-\infty$ when $\omega^{-1}\to 1$.

Therefore, we conclude that the behavior of the deep-inelastic structure functions *in the Bjorken limit* for ω^{-1} near 1 can be substantially modified for mesons by confinement effects, and a violation of the Callan-Gross relation for mesons in the Bjorken limit could also be a consequence of the mentioned effects.

We acknowledge discussions with J. L. Cortés, J. Sánchez Guillén, and especially those with A. Cruz whose critical reading of the manuscript and collaboration at the beginning of this work has been most helpful.

[1] P. V. Landshoff, J. C. Polkinghorne, and R. D. Short, Nucl. Phys. B28, 225 (1971).

[2] We are not interested in the breaking of Bjorken scaling, which has been already investigated [see J. C. Polkinghorne, Nucl. Phys. B108, 253 (1976)] showing

that a consistent covariant parton model can be formulated.

[3]V. Azcoiti, J. L. Alonso, and A. Cruz, Nucl. Phys. B167, 525 (1980).

[4]R. D. Field and R. P. Feynman, Phys. Rev. D 15, 2590 (1977).

[5]We consider for simplicity quarks of spin 0. The generalization to spin-$\frac{1}{2}$ quarks is straightforward and can be seen in Refs. 1 and 3.

[6]It must be noted that if any amplitude involving quarks and at *least two hadrons* is zero on the mass shell, then, because of unitarity and crossing, any on-mass-shell amplitude involving quarks and *only one hadron* is also zero.

[7]The contribution of T_+ to νW_2 is analogous to the T_- contribtuion.

Structure of the Vacuum and Neutron and Neutrino Oscillations

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In a theory where the fermion numbers B and L are not conserved, the true ground state is a condensate of fermion pairs and antifermion pairs. Because of this, the usual analysis of fermion oscillation based on complete analogy with K-$\overline{K}$ mixing must be revised.

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In the context of grand unified theories, where quarks and leptons are put in the same multiplet, it is easy nowadays to accept B and separately L nonconservation.[1] Even $B-L$ may not be sacrosanct and as we have shown in an earlier paper, a nonminimal SU(5) can directly lead to neutron oscillation.[2] The surprising thing about neutron oscillation is that, contrary to first expectations, the *free* neutron oscillation time can be as short as $\sim 10^7$ sec and may in fact be observable.[2,3]

In theories where such fermion numbers as B and L are not conserved, the structure of the vacuum is very much richer than the simple bare Fock vacuum. The new physical vacuum is a condensed state of fermion pairs and antifermion pairs. This leads to new physical effects which are observable in oscillation experiments. Therefore the usual expectation that neutron and neutrino oscillation phenomenology is completely analogous to K-$\overline{K}$ mixing[4] must be revised.

For neutron oscillations, the numerical effect of this new fermion pairing is probably too small to be detectable. For neutrino oscillation,[5] where the neutrinos have *vanishing* Dirac mass, the usual expectation is that ν_e will oscillate only into ν_μ, ν_τ. Our conclusion is that ν_e can also oscillate into $\overline{\nu}_e, \overline{\nu}_\mu, \overline{\nu}_\tau$. Further, the probability amplitudes for ν_e into ν_μ, ν_e, etc., are themselves also modified.

It is instructive to begin our discussion with a simplified fermion-number–nonconserving Lagrangian[6]

$$\mathcal{L} = -\tfrac{1}{2}\overline{\psi}\gamma\cdot\overleftrightarrow{\partial}\psi - m\overline{\psi}\psi - \tfrac{1}{2}\xi\overline{\psi}C^{\dagger}\psi + \tfrac{1}{2}\xi^*\overline{\psi}C\overline{\overline{\psi}}, \qquad (1)$$

where m is the Dirac mass term and ξ the new Majorana mass term. In terms of the usual a,b operators in the Fourier decomposition of ψ, we find the Hamiltonian to be given by $[\omega \equiv +(p^2+m^2)^{1/2}]$

$$H = \int \frac{d^3p}{(2\pi)^3\omega}\left[\omega(a_p{}^{\dagger}a_p + b_p{}^{\dagger}b_p) - \frac{m}{\omega}(\xi\, b_p{}^{\dagger}a_p + \xi^* a_p{}^{\dagger}b_p) + \frac{\xi}{2\omega}\tilde{a}_p\sigma_2\vec{\sigma}\cdot\vec{\mathrm{p}}\,a_{-p} + \mathrm{H.c.} + \frac{\xi^*}{2\omega}\tilde{b}_p\sigma_2\vec{\sigma}\cdot\vec{\mathrm{p}}\,b_{-p} + \mathrm{H.c.}\right]. \qquad (2)$$

Equation (2) clearly exhibits, besides the fermion-antifermion transition term, the new pairing interactions. As is well known from solid state and nuclear physics, these pairing terms cause a condensation of fermion pairs and antifermion pairs in the ground state. This new vacuum is orthogonal to the Fock vacuum.

The structure of this vacuum may be obtained by diagonalizing (2). The diagonal form of H is given

