down, there appearing to be more $h^+$ than $h^-$ from the $n$ target also. These changes in the hadron composition occur in the kinematic region $0.4 < x < 0.85$, a region populated by the decay products or fragments of the $\gamma^*$ in any diffractive model of $\gamma^*$-nucleon interactions. Since the $\gamma^*$ is neutral, the charge asymmetries in electroproduction make any such diffractive model less attractive than in nearly symmetric photoproduction. The above changes from the charge- and isospin-symmetric hadrons of photoproduction to the asymmetric hadrons of electroproduction take place in the $q^2$ range in which scaling begins.[2]

We wish to point out that the behavior shown in Fig. 2 has a natural explanation in a quark-parton model. In such a model the $\gamma^*$ strikes $p$-type (charge $+\frac{2}{3}$) valence quarks in preference to $n$-type (charge $-\frac{1}{3}$) valence quarks. These struck quarks fragment in the $\gamma^*$ fragmentation region, the $p$ type preferentially to $\pi^+$, the $n$ type preferentially to $\pi^-$. This gives a net $\pi^+$ excess for the proton, and a smaller $\pi^+$ excess for the neutron. This model gives a testable prediction[11] for the pion multiplicities in our $x$ range:

$$R=\frac{\int_1^\infty(N_n^+-N_n^-)F_1^n(\omega)\,d\omega/\omega^2}{\int_1^\infty(N_p^+-N_p^-)F_1^p(\omega)\,d\omega/\omega^2}=\frac{2}{7}=0.29. \qquad (7)$$

Here $\omega$ is the scaling variable $(q^2+M^2-s)/q^2$, $F_1(\omega)$ is a known inelastic structure function, and $p$ and $n$ represent the proton and neutron. We are able to test this prediction with our data only over the limited range $3<\omega<60$, and here compute the value $R=0.24\pm0.28$. Clearly a more precise test of relation (7) is needed. A more detailed discussion of these results in relation to the quark-parton model is reported separately.[12]

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†Present address: Rutherford High Energy Laboratory, Chilton, Didcot, Berkshire, England.

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# Motion of Neutrinos in Charged Matter

L. F. Landovitz
*Belfer Graduate School of Science, Yeshiva University, New York, New York 10033*

and

W. Schreiber
*Staten Island Community College, Staten Island, New York 10301*

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The motion of neutrinos in the presence of charged matter is analyzed. Implications for stellar neutrino emission are considered.

In this paper, we shall consider the interaction of a two-component neutrino with charged matter. A two-component neutrino does not have an interaction with the electromagnetic field.[1] However, there is an interaction with the source of the electromagnetic field, i.e., charges and currents. We shall also explore, albeit briefly, some possible astrophysical implications of this type of interaction.

The interaction of a neutrino with an electromagnetic field is expressible by the vertex function[2]

$$\Lambda = F(q^2)\bar{u}(p_\nu')[\tfrac{1}{2}(1-\gamma_5)]\gamma_\mu[\tfrac{1}{2}(1+\gamma_5)]u(p_\nu)A_\mu(q), \tag{1}$$

where

$$F(q^2) \cong F'(0)q^2. \tag{2}$$

The expression $F(q^2)$ is of order $eG$, where $e$ is the electromagnetic coupling constant and $G$ is the Fermi weak-coupling constant. In coordinate space, the quantity $q^2A_\mu(q)$ which appears in Eq. (1) becomes

$$\Box^2 A_\mu(x) = J_\mu(x). \tag{3}$$

To lowest order, the $S$ matrix is of the form $\langle f|H|i\rangle$.

Guided by the form of Eq. (1), we express the motion of a two-component neutrino in charged currents $J$ by the equation

$$i\,\partial\psi/\partial t = \vec{\sigma}\cdot\vec{p}\psi + g\vec{\sigma}\cdot\vec{J}\psi + g\rho\psi, \tag{4}$$

$$J = (\vec{J}, i\rho). \tag{5}$$

We consider the motion of neutrinos in charged distributions rotating with angular velocity $\Omega$. The current $\vec{J}$ is expressed by

$$\vec{J} = \rho\vec{\sigma}\cdot(\vec{\Omega}\times\vec{x}). \tag{6}$$

We choose the axis of rotation to be the $z$ axis and define

$$\epsilon = E - g\rho. \tag{7}$$

Equation (4) now becomes

$$[-\vec{\sigma}\cdot\vec{p} + g\rho\vec{\sigma}\cdot(\vec{\Omega}\times\vec{x})]\psi = \epsilon\psi. \tag{8}$$

One solution to Eq. (8) is obtained by setting

$$\psi = \begin{pmatrix} a_+(r)e^{i(m-1/2)\theta} \\ a_-(r)e^{i(m-1/2)\theta} \end{pmatrix}\exp(ip_3z), \tag{9}$$

where $m$ is the $z$ component of the total angular momentum of the neutrino. The resulting expression for the energy levels is

$$\epsilon = \{p_3^{\,2} + 4\omega[n + \tfrac{1}{2}|m - \tfrac{1}{2}| + \tfrac{1}{2}(m - \tfrac{1}{2}) + 1]\}^{1/2}, \quad n = 0, 1, \dots, \tag{10}$$

where

$$\omega = g\rho\Omega. \tag{11}$$

If the canonical transformation

$$\psi \to \exp(-i\omega x_1x_2)\psi, \tag{12}$$

is performed, the solution obtained is

$$\psi = \begin{pmatrix} u_+(x_1) \\ u_-(x_1) \end{pmatrix}\exp(ip_2x_2)\exp(ip_3x_3). \tag{13}$$

The resulting energy levels are

$$\epsilon = [p_3^{\,2} + 4\omega(n+1)]^{1/2}, \quad n = 0, 1, \dots. \tag{14}$$

A charged particle of mass $\mu$ moving in a uniform magnetic field has similar solutions. Here, too, one has a solution involving Laguerre polynomials having energy levels

$$E = \{p_3^2 + \mu^2 + 4\omega[n + \tfrac{1}{2}|m| + \tfrac{1}{2}m + 1]\}^{1/2}, \tag{15}$$

$$\omega = eB/2, \tag{16}$$

and a solution involving Hermite polynomials having energy levels

$$E = [p_3^2 + \mu^2 + 4\omega(n+1)]^{1/2}. \tag{17}$$

The similarity between the charged-particle motion and the neutrino motion implies that a neutrino packet emitted from the core of a star will have a mean square radius associated with it. Specifically, one obtains a mean square radius

$$R \sim p/\omega. \tag{18}$$

If $\omega$ is sufficiently large, e.g., the density or rotational velocity extremely high, or if there were a new interaction which would increase the coupling constant several orders of magnitude, then neutrinos would be trapped inside the matter and never emerge.

If $\omega^{-1/2} \lesssim R_\odot^{\,3}$ ($\omega$ is very model dependent, and thus difficult to estimate), the energy levels are approximately of the form

$$\epsilon = [p_3^2 + 4\omega(n+1)]^{1/2}, \tag{19}$$

rather than the free-neutrino form

$$E = (p_1^2 + p_2^2 + p_3^2)^{1/2}, \tag{20}$$

which has implications with respect to the solar neutrino problem.

Recent experiments indicate a considerably smaller counting of neutrinos from the sun than was expected.[4] We have investigated the respective counting rates implied by Eqs. (19) and (20) for the process $\nu + {}^{37}\mathrm{Cl} \rightarrow e^- + {}^{37}\mathrm{Ar}$ of the experiment from phase-space considerations. Assuming the same flux at the earth for both cases, the relevant quantity is

$$R = R_2/R_1, \tag{21}$$

where

$$R_i = \int_{\Delta M + m_e}^{\Delta M' - m_e} dp\, \rho_i(p)(p - \Delta M)[(p-\Delta M)^2 - m_e^2]^{1/2} F(Z,p) \left[\int_0^{\Delta M' - m_e} \rho_i(p)\, dp\right]^{-1}. \tag{22}$$

In the above equation, $\Delta M$ and $\Delta M'$ are the absolute values of the ${}^{37}\mathrm{Cl}$-${}^{37}\mathrm{Ar}$ mass difference and the solar neutrino producing ${}^{8}\mathrm{B}$-${}^{8}\mathrm{Be}^*$ mass difference, respectively, and $F(Z,p)$ is the Coulomb correction factor. For the energy form of Eq. (20), one obtains

$$\rho_1(p) \propto p^2(\Delta M' - p)[(\Delta M' - p)^2 - m_e^2]^{1/2} F(Z', p), \tag{23}$$

while for the energy form of Eq. (19), one obtains

$$\rho_2(p) \propto (\Delta M' - p)[(\Delta M' - p)^2 - m_e^2]^{1/2} F(Z', p). \tag{24}$$

The different phase-space factor occurs because the form of Eq. (19) rather than Eq. (20) results in the replacement

$$\int d^3p_\nu \rightarrow \sum\nolimits_n \int d\rho_3. \tag{25}$$

The sum is replaced by an integral using the first term of the Euler-Maclaurin summation formula which should be very accurate. The value obtained for $R$ is about $\frac{1}{3}$. Thus, from phase-space considerations alone, the neutrino counting rate may be reduced by a factor of 3.

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# Vacuum Polarization and the Quark-Parton Puzzle

A. Casher,* J. Kogut, and Leonard Susskind
*Tel-Aviv University, Tel-Aviv, Israel*
(Received 2 July 1973)

We suggest a solution to the question of why isolated quark-partons are not observed. It is argued that gauge theories can have the light-cone behavior of quark fields without quark production in deep-inelastic reactions. Our ideas are illustrated in the exact solution of quantum electrodynamics in one space and one time dimension. The theory's crucial physical property is the extreme polarizability of its vacuum. Predictions for deep inelastic final states are presented.

The quark-parton model has provided a simple intuitive guide for predicting and interpreting many features of deep-inelastic scattering processes.[1] Its success, however, underscores an old problem which was born when the quark model was invented: Why are isolated quarks not observed? In the context of deep-inelastic scattering this familiar problem becomes even more puzzling. The reason for this can be found in the fact that the parton model is based on soft (superrenormalizable or cutoff) field theories in which partons experience just soft, finite forces.[2] This suggests that when a parton absorbs a very virtual photon of energy $\nu$, it should propagate essentially freely for lab-frame distances which grow proportionally with $\nu$. Once this distance exceeds the size of a hadron (1 fm, say), one might naively expect that quarks would then be produced. In this Letter, however, we shall argue that this is not necessarily the case. If, in fact, the underlying constituent field theory has a sufficiently polarizable vacuum, quark-partons will never be produced.

Consider first an exactly soluble example of this phenomenon: quantum electrodynamics in one space and one time dimension. We shall argue that the deep-inelastic structure functions of this theory have the scaling laws of underlying free Fermi fields, although only massive bosons appear as asymptotic stable particles. The analog of triality in this simplified model is fermion number. The complete solution of this theory was first given by Schwinger.[3] The exact Green's functions can be obtained through a formal solution of the theory's equations of motion. We shall sketch the derivation here and leave the details and discussion to a lengthier exposition to appear elsewhere.[4] The familiar Lagrangian reads

$$\mathcal{L} = \bar{\psi} i\gamma^{\mu}\partial_{\mu}\psi - \tfrac{1}{4}F^{\mu\nu}F_{\mu\nu} - e\bar{\psi}\gamma^{\mu}\psi A_{\mu}, \tag{1}$$

and the equations of motion are

$$\begin{gathered} \gamma^{\mu}(i\partial_{\mu} - eA_{\mu})\psi = 0, \\ j^{\mu} \equiv e\bar{\psi}\gamma^{\mu}\psi = \partial_{\nu}F^{\nu\mu}, \\ F^{\mu\nu} = \partial^{\mu}A^{\nu} - \partial^{\nu}A^{\mu}. \end{gathered} \tag{2}$$

A careful definition of the current leads to the equation[3]

$$(\Box + m^2)j^{\mu} = 0, \tag{3}$$

where $m^2 = e^2/\pi$. Thus, the spectrum of the theory contains free massive bosons. It is useful to introduce a scalar function $\varphi(x)$ such that

$$j^{\mu} = \epsilon^{\mu\nu}\partial_{\nu}\varphi, \quad \varphi(z,t) = \int_{-\infty}^{z} j^0(z',t)\,dz'; \tag{4}$$

$\varphi$ will later be interpreted as a dipole density. $\varphi$ satisfies the equation of motion and canonical commutation relation

$$\begin{gathered} (\Box + m^2)\varphi = 0, \\ [\varphi(z,t),\ \dot{\varphi}(z',t)] = im^2\delta(z-z'). \end{gathered} \tag{5}$$

The equation of motion for $A_{\mu}$ in the Lorentz gauge follows from Eq. (2),

$$\Box A_{\mu} = j_{\mu}. \tag{6}$$