Density matrix of relativistic nuclear matter

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We investigate the structure of the density matrix of relativistic nuclear matter at large momentum ($k \gg k_F$) making use of the $\sigma+\omega$ model of Walecka. The density matrix may be written in terms of spinors describing nucleons of modified mass and also has exotic components that contain spinors describing negative-energy states ("antinucleons"). We calculate the modification of the mean-field density matrix due to the admixture of two-particle, two-hole states in the ground state of the relativistic theory and find that the excitation of the exotic components of the density matrix is so strong and *coherent* as to preclude the use of perturbation theory. We suggest that an expansion in terms of reaction matrices will not improve the situation and conclude that a consistent picture may be obtained only if one limits oneself to a space spanned by positive-energy spinors (describing nucleons of shifted mass). We also consider the exchange of pions in the calculation of high-momentum components of the density matrix. However, in the case of pion exchange, we find that sensible results may not be obtained (even for the nonexotic components) unless we include tensor and short-range correlations. (The latter calculations have not been performed as yet.) Further, we note that a theory without vertex cutoffs (meson-nucleon form factors) leads to totally unacceptable results, as the depletion of the Fermi sea is again much too large.

I. INTRODUCTION

There has been an ongoing interest in the study of high-momentum components in nuclei and the associated depletion of occupied states due to short-range and tensor correlations.[1–4] High-momentum components appear if one goes beyond the mean-field approximation and includes two-particle, two-hole excitations in the ground state. Studies of such components may be made in the case of nuclear matter, if one uses Brueckner theory. That theory is usually applied within the context of Schrödinger dynamics. The question may be raised, however, as to the validity of the Schrödinger theory when one considers momenta which are comparable to the nucleon mass. For example, in the study of y scaling[5] one may attempt to obtain information concerning momentum components as large as 0.8 GeV/c. In that regime we believe a fully relativistic description of the system is necessary.

Relativistic theories of nuclear structure have been introduced in recent years.[6,7] However, the formalism is still under development, since there is no general agreement as to the role of negative-energy states in such theories. On the other hand, advocates of the use of the relativistic formalism do agree that there is good evidence for the presence of very large (Lorentz) scalar and vector fields in the nuclear medium. The success of the relativistic formalism in describing various experimental data has its origin in the action of these large fields which serve to shift the *mass* in the Dirac equation away from the value it would have in vacuum, while not producing major changes in the *energy* of the nucleon.[7]

Various calculations have been made using the relativistic models of nuclear structure.[8] However, the density matrix of relativistic nuclear matter has not been investigated in a systematic fashion. In the relativistic Brueckner-Hartree-Fock theory, such calculations are quite complex. Therefore, in this work we will limit our considerations to the Walecka model[6] and perform some exploratory calculations in the asymptotic region $k \gg k_F$.

We note that the density matrix of relativistic nuclear matter is a 4×4 Dirac matrix $\rho_{\alpha\beta}(k)$, in contrast to the density matrix of the Schrödinger theory, where the nucleons are described by two-component spinors. We suggest that if one probes nuclear matter at high momentum and at low-energy transfer, the response will be governed by $\rho_{\alpha\beta}(k)$, to the extent that the impulse approximation may be used to describe this response.

We may recall that in the theory of relativistic nuclear matter outlined in Ref. 7, we introduced the positive- and negative-energy solutions of the Dirac equation, $f(\mathbf{p},s)$ and $h(\mathbf{p},s)$, respectively. The spinor $f(\mathbf{p},s)$ was proportional to the free spinor $u(\mathbf{p},s,\tilde{m})$, which was parametrized by the shifted mass variable:

$$\tilde{m}(\mathbf{p})=m_N+\tfrac{1}{4}\mathrm{Tr}\Sigma(\mathbf{p}) \ . \tag{1.1}$$

Here $\Sigma(\mathbf{p})$ is the self-energy operator which appears in the Dirac equation. In the relativistic Brueckner-Hartree-Fock theory, the energy of the system was expressed only in terms of the $f(\mathbf{p},s)$, which describe positive-energy *nucleon* quasiparticle degrees of freedom. On the other hand, if one does not truncate the theory in that manner, the negative-energy states play a role, and in particular, the density matrix, which we here consider for $k \gg k_F$, will have the structure

$$\rho_{\alpha\beta}(\mathbf{k})=\sum_{ss'}[f_\alpha(\mathbf{k},s)D^{++}_{ss'}(\mathbf{k})\bar{f}_\beta(\mathbf{k},s') +f_\alpha(\mathbf{k},s)D^{+-}_{ss'}(\mathbf{k})\bar{h}_\beta(\mathbf{k},s') +h_\alpha(\mathbf{k},s)D^{-+}_{ss'}(\mathbf{k})\bar{f}_\beta(\mathbf{k},s') +h_\alpha(\mathbf{k},s)D^{--}_{ss'}(\mathbf{k})h_\beta(\mathbf{k},s')] . \quad (1.2)$$

[If we do not make the Dirac indices explicit, we will use the notation $f_s(\mathbf{p})=f(\mathbf{p},s)$, $u_s(\mathbf{p})=u(\mathbf{p},s)$, etc. For simplicity, we have suppressed reference to isospin.] A complete calculation of this density matrix is a lengthy project, and we are here concerned with making some asymptotic estimates for $k \gg k_F$. From our past experience,[7] we know that $D^{++}_{ss'}(k)$ will be sensitive to short-range correlations; however, reasonable results for various observables may be obtained by adjusting the coupling constants of the sigma and omega fields. (These adjusted constants may be considered to contain some of the effects of short-range correlations.) On the other hand, a quantity such as $D^{--}(k)$ was found to be remarkably insensitive to correlation effects in calculations we have performed previously.[7] Therefore, in the calculations reported here, which we consider to be exploratory in nature, we will adopt the philosophy of the Walecka model and use adjustable parameters, rather than attempt a detailed calculation of short-range correlation effects. From this point of view, we are therefore investigating whether the Walecka model provides a description which is stable against admixtures of two-particle, two-hole states into the ground state. As we will see, even with the use of form factors at the vertices, exotic components such as $D^{--}(k)$ are so large as to indicate nonconvergence in this method of calculation.

We recall that it has been suggested that antinucleon excitations will be suppressed by form factors having their origin in the composite nature of the nucleons and antinucleons.[9] Therefore, in a phenomenonological description of relativistic nuclear matter, we may wish to keep only the first term of Eq. (1.2). That suggestion is in accord with the procedures adopted in the relativistic Brueckner-Hartree-Fock (RBHF) theory we have developed previously.[7]

The density matrix of nuclear matter has recently been discussed by Jaminon and Mahaux[10] in the case that the relativistic Brueckner-Hartree-Fock approximation is used. The method adopted in Ref. 10 is to write the nucleon potential as a sum of two terms, one of which is frequency independent and corresponds to what would be obtained in a Hartree-Fock approximation. The other (dispersive) term satisfies a dispersion relation relating the real and imaginary parts. The imaginary parts of the scalar and vector potentials are taken from Ref. 11 and are used to determine the real parts by making use of dispersion relations. Occupation probabilities for hole and particle states are then obtained by differentiation with respect to energy of the potentials obtained in this manner. Depletion of the occupied states is found to be about 5% on average, which is a significantly smaller depletion than that found in nonrelativistic Brueckner-Hartree-Fock theory.[1–4] We should note, however, that the work of Ref. 10 deals only with the part of the density matrix of Eq. (1.2) specified by $D^{++}(\mathbf{k})$. In the work reported here, we are mainly concerned with making some asymptotic estimates of the other components of the density matrix specified by the values of D^{+-}, D^{-+}, and D^{--}.

The plan of our work is as follows. In Sec. II we discuss the structure of the density matrix and describe a calculation in the asymptotic domain $k \gg k_F$. In Sec. III we present the results of our calculations and some further discussion.

II. DENSITY MATRIX OF RELATIVISTIC NUCLEAR MATTER

The positive- and negative-energy solutions of the Dirac equation for nucleons embedded in infinite nuclear matter can be represented by the spinors[7]

$$f_s(\mathbf{k})=\left[\frac{E(\mathbf{k})\tilde{m}}{\tilde{E}(\mathbf{k})m_N}\right]^{1/2}u_s(\mathbf{k}) , \quad (2.1)$$

$$h_s(\mathbf{k})=\left[\frac{E(\mathbf{k})\tilde{m}}{\tilde{E}(\mathbf{k})m_N}\right]^{1/2}v_{-s}(-\mathbf{k}) , \quad (2.2)$$

where $u_s(\mathbf{k})$ and $v_s(\mathbf{k})$ are free Dirac spinors with mass parameter $\tilde{m}$; $E(\mathbf{k})=(m_N^2+\mathbf{k}^2)^{1/2}$ and $\tilde{E}(\mathbf{k})=(\tilde{m}+\mathbf{k}^2)^{1/2}$. Here m_N is the mass of a free nucleon and $\tilde{m}$ is the mass in nuclear matter. The Green's function for a nucleon propagating in nuclear matter can be expressed in terms of the positive- and negative-energy spinors of Eqs. (2.1) and (2.2):[7]

$$G(k)=G^+(k)+G^-(k) \quad (2.3)$$

$$=\left[\frac{m_N}{E(\mathbf{k})}\right]^{1/2}\sum_s\left[\frac{f_s(\mathbf{k})\bar{f}_s(\mathbf{k})}{k^0-\epsilon^+(\mathbf{k})+i\eta}+\frac{h_s(\mathbf{k})\bar{h}_s(\mathbf{k})}{k^0-\epsilon^-(\mathbf{k})-i\eta}\right] . \quad (2.4)$$

In nuclear matter $\epsilon^\pm(\mathbf{k})=B\pm\tilde{E}(\mathbf{k})$ and $\tilde{m}=m_N+A$, where $B\approx 0.3$ GeV and $A\approx -0.4$ GeV.[6,7]

We now consider the correction to the nuclear matter density matrix due to excitation of the nucleons out of the Fermi sea via meson exchange (see Fig. 1). The mesons considered here are σ, ω, and π, although one can also add the contribution of the ρ. The nuclear density matrix is then given by

$$\rho(k)=\rho_0(k)+\delta\rho(k) , \quad (2.5)$$

where $\rho_0(k)$ is the contribution to the density matrix of the particles in the Fermi sea. In the absence of depletion, we would have

$$\rho_0(k)=\sum_{s,s'}f_s(\mathbf{k})\bar{f}_{s'}(\mathbf{k})\Theta(k_F-|\mathbf{k}|)\delta[k^0-\epsilon^+(\mathbf{k})]\delta_{ss'} . \quad (2.6)$$

The Fermi momentum is taken to be $k_F=0.272$ GeV. $\delta\rho(k)$ is the correction due to nucleon excitations and includes contributions containing negative-energy spinors, if we use the full propagator of Eq. (2.4):

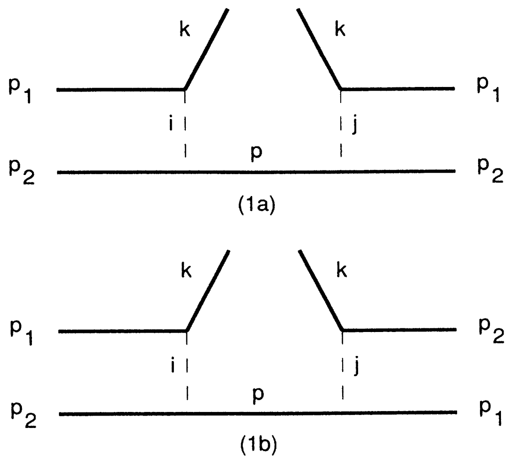

FIG. 1. Diagrams for one-meson-exchange contributions to the nuclear density matrix: (a) direct term and (b) exchange term. Here p_1 and p_2 denote nucleons in the Fermi sea, while $|\mathbf{p}|>k_F$. Such diagrams appear as parts of larger Feynman diagrams, when one calculates the electromagnetic response tensor of nuclear matter, for example.

$$\delta\rho(k)=\delta\rho^{++}(k)+\delta\rho^{+-}(k)+\delta\rho^{-+}(k)+\delta\rho^{--}(k)\ . \tag{2.7}$$

We will discuss $\delta\rho(k)$ for $|\mathbf{k}|\gg k_F$. The various terms comprising $\delta\rho(k)$ are defined in terms of a density matrix $d(k)$, which does not include the external fermion lines of the diagrams of Fig. 1:

$$\delta\rho^{++}(k)=G^{+}(k)d(k)G^{+}(k)\ , \tag{2.8}$$

$$\delta\rho^{+-}(k)=G^{+}(k)d(k)G^{-}(k)\ , \tag{2.9}$$

$$\delta\rho^{-+}(k)=G^{-}(k)d(k)G^{+}(k)\ , \tag{2.10}$$

$$\delta\rho^{--}(k)=G^{-}(k)d(k)G^{-}(k)\ . \tag{2.11}$$

[Note that the quantities appearing in Eqs. (2.9)–(2.11) are particular to the Walecka model,[6] since in the relativistic Brueckner-Hartree-Fock theory[7] we dropped the second term of the propagator in Eq. (2.4) and developed the theory for only positive-energy quasiparticles of modified mass $\tilde{m}$.]

In the case of identical nucleons, one has to include, in addition to the direct term shown in Fig. 1(a), the exchange diagram of Fig. 1(b). The use of Eq. (2.4) for the Green's function leads to a specification of separate nucleon and antinucleon contributions. For simplicity, we make the assumption that the form factors for meson-nucleon and meson-antinucleon coupling are equal. This might be a reasonable assumption; however, there are objections to its validity.[9]

Note that, for symmetric nuclear matter, both $\rho_0(k)$ and $\delta\rho(k)$ are independent of the nucleon isospin. Thus the density matrix $d(k)$ can be written in a compact form. Since the interaction shown in Fig. 1 is a sum of meson exchanges, the contribution to $\delta\rho(k)$ contains various products which may be labeled by indices which refer to the character of the exchanged mesons. With $k=p_1+p_2-p$ (see Fig. 1), we have

$$\begin{aligned} d(k)=\frac{m_N^3}{(2\pi)^5}\int\frac{d^4p_1d^4p_2}{E_1E_2}\delta(p_1^0-\epsilon_1^+)\delta(p_2^0-\epsilon_2^+)\frac{\delta(p^0-\epsilon^+)}{E}\\ \times\Theta(k_F-|\mathbf{p}_1|)\Theta(k_F-|\mathbf{p}_2|)\Theta(|\mathbf{p}|-k_F)\sum_{i,j}G^i(Q^2)G^j(Q^2)(L_D^{ij}-L_E^{ij})\ . \end{aligned} \tag{2.12}$$

Here the sum is over i and j, which denote σ, ω, and π meson contributions. Thus there are σ-σ, σ-ω, ω-ω, σ-ω, etc., contributions in Eq. (2.12). Further, $G^i(Q^2)$ denotes the product of the propagator of meson i times the meson-nucleon (or antinucleon) form factors which appear at the vertices. L_D^{ij} and L_E^{ij} are operators which arise from the direct and exchange terms and refer to exchange of mesons i and j. For example,

$$\begin{aligned} L_D^{ij}&=T_D^{ij}\sum_{s_1s_2s}\Gamma^i f_{s_1}(\mathbf{p}_1)\bar{f}_s(\mathbf{p})\Gamma^i f_{s_2}(\mathbf{p}_2)\bar{f}_{s_2}(\mathbf{p}_2)\Gamma^j f_s(\mathbf{p})\bar{f}_{s_1}(\mathbf{p}_1)\Gamma^j\\ &=T_D^{ij}\frac{E_1E_2E}{\tilde{E}_1\tilde{E}_2\tilde{E}}\left[\frac{\tilde{m}}{m_N}\right]^3\Gamma^i\frac{\tilde{m}+\tilde{\not{p}}_1}{2\tilde{m}}\Gamma^j\mathrm{Tr}\left[\Gamma^j\frac{\tilde{m}+\tilde{\not{p}}}{2\tilde{m}}\Gamma^i\frac{\tilde{m}+\tilde{\not{p}}_2}{2\tilde{m}}\right]\ . \end{aligned} \tag{2.13}$$

Here T_D^{ij} is the isospin factor of the direct term:

$$T_D^{ij}=\sum_{t_1t_2t}\langle t_k|T^i|t_1\rangle\langle t_1|T^j|t_k\rangle\langle t|T^i|t_2\rangle\langle t_2|T^j|t\rangle\ . \tag{2.14}$$

Further,

$$L_E^{ij}=T_E^{ij}\sum_{s_1s_2s}\Gamma^i f_{s_1}(\mathbf{p}_1)\bar{f}_s(\mathbf{p})\Gamma^i f_{s_2}(\mathbf{p}_2)\bar{f}_{s_1}(\mathbf{p}_1)\Gamma^j f_s(\mathbf{p})\bar{f}_{s_2}(\mathbf{p}_2)\Gamma^j \tag{2.15}$$

$$=T_E^{ij}\frac{E_1E_2E}{\tilde{E}_1\tilde{E}_2\tilde{E}}\left[\frac{\tilde{m}}{m_N}\right]^3\Gamma^i\frac{\tilde{m}+\tilde{\not{p}}_1}{2\tilde{m}}\Gamma^j\frac{\tilde{m}+\tilde{\not{p}}}{2\tilde{m}}\Gamma^i\frac{\tilde{m}+\tilde{\not{p}}_2}{2\tilde{m}}\Gamma^j\ , \tag{2.16}$$

with the isospin factor

$$T_E^{ij}=\sum_{t_1t_2t}\langle t_k|T^i|t_1\rangle\langle t_1|T^j|t_k\rangle\langle t|T^i|t_2\rangle\langle t_2|T^j|t\rangle\delta_{t_1t_2}\ . \tag{2.17}$$

In these expressions Γ^i is the vertex operator for the meson i and T^i is the associated isospin operator. These operators may be specified, if we write the following meson-nucleon interaction Lagrangian:

$$\mathcal{L}_I=-g_\sigma\bar{\Psi}_N\sigma\Psi_N-g_\omega\bar{\Psi}_N\omega_\mu\gamma^\mu\Psi_N-\frac{f_\pi}{m_\pi}\bar{\Psi}_N\gamma_5\gamma_\mu\boldsymbol{\tau}\cdot\partial^\mu\boldsymbol{\pi}\Psi_N\ . \tag{2.18}$$

Since σ and ω are isosinglets, $T^\sigma=T^\omega=1$. For the pion, $T^\pi i=\tau_i$. The vertex operators are $\Gamma_\sigma=1$, $\Gamma_\mu^\omega=\gamma_\mu$, and $\Gamma^\pi=\gamma_5 \not{q}$, where q_μ is the momentum transfer. The coupling constants are absorbed in the definition of the meson-nucleon form factor. Using the energy-conserving delta functions, we perform the integration over the energies of the nucleons labeled by p_1 and p_2. We have

$$d(k)=\frac{1}{(2\pi)^5}\int\frac{d^3p_1d^3p_2}{4\tilde{E}_1\tilde{E}_2}\,\frac{\delta(p^0-\epsilon^+)}{2\tilde{E}}\Theta(k_F-|\mathbf{p}_1|)\Theta(k_F-|\mathbf{p}_2|)\Theta(|\mathbf{p}|-k_F)\sum_{i,j}G^i(Q^2)G^j(Q^2)(I_D^{ij}-I_E^{ij})\ , \tag{2.19}$$

with

$$I_D^{ij}=T_D^{ij}\Gamma^i(\tilde{m}+\tilde{\not{p}}_1)\Gamma^j\mathrm{Tr}[\Gamma^j(\tilde{m}+\tilde{\not{p}})\Gamma^i(\tilde{m}+\tilde{\not{p}}_2)]\ , \tag{2.20}$$

$$I_E^{ij}=T_E^{ij}\Gamma^i(\tilde{m}+\tilde{\not{p}}_1)\Gamma^j(\tilde{m}+\tilde{\not{p}})\Gamma^i(\tilde{m}+\tilde{\not{p}}_2)\Gamma^j\ . \tag{2.21}$$

We proceed to calculate explicit expressions for I_D^{ij} and I_E^{ij}. To avoid a cumbersome notation, we drop the tilde from the nucleon momenta. We obtain

$$I_D^\sigma=4(2)(\tilde{m}^2+p\cdot p_2)(\tilde{m}+\not{p}_1)\ , \tag{2.22}$$

$$I_D^\omega=8(2)[\tilde{m}(2\tilde{m}^2-p\cdot p_2)+p_1\cdot p_2\not{p}_2+p_1\cdot p_2\not{p}-\tilde{m}^2\not{p}_1]\ , \tag{2.23}$$

$$I_D^{\sigma,\omega}=8(2)\tilde{m}[p_1\cdot p+p_1\cdot p_2+\tilde{m}(\not{p}+\not{p}_2)]\ , \tag{2.24}$$

$$I_D^\pi=32(6)\tilde{m}^2(\tilde{m}^2-p\cdot p_2)[(\tilde{m}^2-p_2\cdot p)(\tilde{m}+\not{p}_1)+(p_1\cdot p_2-p_1\cdot p)(\not{p}-\not{p}_2)]\ , \tag{2.25}$$

$$I_D^{\pi,\sigma}=0(0)\ , \tag{2.26}$$

$$I_D^{\pi,\omega}=0(0)\ , \tag{2.27}$$

where the number in the parentheses denotes the isospin factor for the particular process. (Note the absence of cross terms in the case of π exchange.)

Because of the absence of a trace in Eq. (2.21), the operators for the exchange terms are more complicated. However, we can bring them to a simplified form by noting that the remaining part of the integrand is symmetric under the interchange of $p_1\leftrightarrow p_2$, and therefore only the symmetric part of I_E^{ij} contributes to the density matrix. Consequently, we write only the symmetric part of the operators I_E^{ij} of Eq. (2.21):

$$I_E^\sigma=[\tilde{m}(\tilde{m}^2+p_1\cdot p_2+p_1\cdot p+p_2\cdot p)+(\tilde{m}^2+p_1\cdot p)\not{p}_2+(\tilde{m}^2+p_2\cdot p)\not{p}_1+(\tilde{m}^2-p_1\cdot p_2)\not{p}]\ , \tag{2.28}$$

$$I_E^\omega=4[\tilde{m}(p_1\cdot p_2+p_1\cdot p+p_2\cdot p-2\tilde{m}^2)+\tilde{m}^2(\not{p}_1+\not{p}_2)+(\tilde{m}^2-2p_1\cdot p_2)\not{p}]\ , \tag{2.29}$$

$$I_E^{\sigma,\omega}=2[\tilde{m}(4\tilde{m}^2+p_1\cdot p+p_2\cdot p-2p_1\cdot p_2)+(\tilde{m}^2+2p_2\cdot p)\not{p}_1+(\tilde{m}^2+2p_1\cdot p)\not{p}_2-2\tilde{m}^2\not{p}]\ , \tag{2.30}$$

$$\begin{aligned}I_E^\pi=24\tilde{m}^2[&\tilde{m}[2p_1\cdot p_2(\tilde{m}^2+p_1\cdot p_2-p_1\cdot p-p_2\cdot p)+(p_1\cdot p)^2+(p_2\cdot p)^2-\tilde{m}^4-(p_1\cdot p_2)^2]\\&+[(\tilde{m}^2+p_1\cdot p_2)(\tilde{m}^2+p_1\cdot p_2-p_1\cdot p-p_2\cdot p)+2(p_1\cdot pp_2\cdot p-\tilde{m}^2p_1\cdot p_2)]\not{p}\\&+[(\tilde{m}^2-p_1\cdot p_2)(p_2\cdot p-\tilde{m}^2)+(p_1\cdot p-p_2\cdot p)(\tilde{m}^2+p_2\cdot p)]\not{p}_1\\&+[(\tilde{m}^2-p_1\cdot p_2)(p_1\cdot p-\tilde{m}^2)+(p_2\cdot p-p_1\cdot p)(\tilde{m}^2+p_1\cdot p)]\not{p}_2]\ ,\end{aligned} \tag{2.31}$$

$$\begin{aligned}I_E^{\pi,\sigma}=4\tilde{m}[&2(\tilde{m}^4-p_1\cdot pp_2\cdot p)+\tilde{m}(2p_1\cdot p_2-p_2\cdot p-p_1\cdot p)\not{p}\\&+\tilde{m}(2\tilde{m}^2-p_1\cdot p-p_1\cdot p_2)\not{p}_1+\tilde{m}(2\tilde{m}^2-p_1\cdot p_2-p_2\cdot p)\not{p}_2]\ ,\end{aligned} \tag{2.32}$$

$$\begin{aligned}I_E^{\pi,\omega}=8\tilde{m}[&2(\tilde{m}^2-p_1\cdot p)(\tilde{m}^2-p_2\cdot p)+p_1\cdot p_2(2\tilde{m}^2-p_1\cdot p-p_2\cdot p)\\&+2\tilde{m}\not{p}(\tilde{m}^2+p_1\cdot p_2-p_1\cdot p-p_2\cdot p)+\tilde{m}(p_1\cdot p-p_1\cdot p_2)\not{p}_1+\tilde{m}(p_2\cdot p-p_1\cdot p_2)\not{p}_2]\ .\end{aligned} \tag{2.33}$$

Since we are interested in an asymptotic approximation to $\delta\rho(k)$, we limit our calculation to the case of small values of $|\mathbf{p}_1|/|\mathbf{k}|$ and $|\mathbf{p}_2|/|\mathbf{k}|$, that is, equivalent to the condition $|\mathbf{k}| \gg k_f$. Since k_F is the upper limit for $|\mathbf{p}_1|$ and $|\mathbf{p}_2|$, our results can be expected to be reasonably accurate for $|\mathbf{k}| > 2k_F \approx \tilde{m}$. In that limit we can approximate the integrals over $\mathbf{p}_1$ and $\mathbf{p}_2$ by placing $\mathbf{p}_1 = \mathbf{p}_2 = 0$ in the integrand. The exchange operators, in that limit, are

$$I_D^\sigma \simeq 8\tilde{m}^2(\tilde{m}+\tilde{E})(1+\gamma_0) , \tag{2.34}$$

$$I_D^\omega \simeq 16\tilde{m}^2[2\tilde{m}-\tilde{E}+(2\tilde{E}-\tilde{m})\gamma_0-\mathbf{p}\cdot\boldsymbol{\gamma}] , \tag{2.35}$$

$$I_D^{\sigma,\omega} \simeq 16\tilde{m}^2[(\tilde{m}+\tilde{E})(1+\gamma_0)-\mathbf{p}\cdot\boldsymbol{\gamma}] , \tag{2.36}$$

$$I_D^\pi \simeq 192\tilde{m}^4(\tilde{m}-\tilde{E})^2(\tilde{m}+\tilde{E}\gamma_0-\mathbf{p}\cdot\boldsymbol{\gamma}) , \tag{2.37}$$

$$I_D^{\pi,\sigma} \simeq 0 , \tag{2.38}$$

$$I_D^{\pi,\omega} \simeq 0 , \tag{2.39}$$

and

$$I_E^\sigma \simeq \tfrac{1}{4} I_D^\sigma , \tag{2.40}$$

$$I_E^\omega \simeq 4\tilde{m}^2[(2\tilde{E}-\tilde{m})+(2\tilde{m}-\tilde{E})\gamma_0+\mathbf{p}\cdot\boldsymbol{\gamma}] , \tag{2.41}$$

$$I_E^{\omega,\sigma} \simeq 4\tilde{m}^2[(\tilde{m}+\tilde{E})(1+\gamma_0)+\mathbf{p}\cdot\boldsymbol{\gamma}] , \tag{2.42}$$

$$I_E^\pi \simeq \tfrac{1}{4} I_D^\pi , \tag{2.43}$$

$$I_E^{\pi,\sigma} \simeq 8\tilde{m}^3(\tilde{m}-\tilde{E})[(\tilde{m}+\tilde{E})(1+\gamma_0)-\mathbf{p}\cdot\boldsymbol{\gamma}] , \tag{2.44}$$

$$I_E^{\pi,\omega} \simeq 16\tilde{m}^3(\tilde{m}-\tilde{E})[2\tilde{m}-\tilde{E}+(2\tilde{E}-\tilde{m})\gamma_0-2\mathbf{p}\cdot\boldsymbol{\gamma}] . \tag{2.45}$$

[Note that in the limit $\mathbf{p}\to 0$ (or equivalently $\mathbf{k}\to 0$), $I_E^{ij} = \frac{1}{4} I_D^{ij}$ for all meson-exchange processes.] We perform the momentum integrations, neglecting $\mathbf{p}_1$ and $\mathbf{p}_2$ in the integrand, and use

$$\int d^3p_1 d^3p_2 \Theta(k_F-|\mathbf{p}_1|)\Theta(k_F-|\mathbf{p}_2|) = (\tfrac{4}{3}\pi k_F^3)^2 . \tag{2.46}$$

Thus the density matrix is given by

$$d(k) = \frac{1}{16\pi^2} W(k) \sum_{i,j} G^i(Q^2) G^j(Q^2)(I_D^{ij}-I_E^{ij}) , \tag{2.47}$$

where $W(k)$ is defined by

$$W(k) = \frac{k_F^6}{9\pi\tilde{m}^2\tilde{E}(k)} \delta[k^0+\epsilon^+(\mathbf{k})-2\epsilon^+(0)]\Theta(|\mathbf{k}|-k_F) . \tag{2.48}$$

Note that, because of the approximation adopted in the above expressions, the density matrix is proportional to a delta function defining k^0 in terms of $|\mathbf{k}|$. In the more general case, when one does not neglect the nucleon momenta $\mathbf{p}_1$ and $\mathbf{p}_2$, the matrix $d(k)$ would be a smooth function of k^0.

We now examine the various parts of the density matrix corresponding to the exchanged mesons. We define the meson potentials V_π, V_σ, and V_ω:

$$V_\sigma(q^2) = \frac{g_\sigma^2}{4\pi} \frac{F_\sigma^2(q^2)}{q^2-m_\sigma^2} , \tag{2.49}$$

$$V_\omega(q^2) = \frac{g_\omega^2}{4\pi} \frac{F_\omega^2(q^2)}{q^2-m_\omega^2} , \tag{2.50}$$

$$V_\pi(q^2) = \frac{f_\pi^2}{4\pi} \frac{F_\pi^2(q^2)}{q^2-m_\pi^2} \left[\frac{\tilde{m}}{m_\pi}\right]^2 , \tag{2.51}$$

and

$$\bar{V}_\pi(q^2) = 2\frac{\tilde{E}-\tilde{m}}{\tilde{m}} V_\pi(q^2) . \tag{2.52}$$

Then the direct and exchange parts of the density matrix can be written in an instructive form:

$$d_D(k) = 8\tilde{m}^2 W(k)\{(1+\gamma_0)(\tilde{m}+\tilde{E})[(V_\sigma-V_\omega)^2+3\bar{V}_\pi^2]+3(1-\gamma_0)(\tilde{m}-\tilde{E})(V_\omega^2+\bar{V}_\pi^2)+2\mathbf{k}\cdot\boldsymbol{\gamma}[V_\omega(V_\omega-V_\sigma)+3\bar{V}_\pi^2]\} , \tag{2.53}$$

and

$$d_E(k) = 2\tilde{m}^2 W(k)\{(1+\gamma_0)(\tilde{m}+\tilde{E})[(V_\sigma-V_\omega)^2+3\bar{V}_\pi^2-2\bar{V}_\pi(V_\sigma-V_\omega)]+3(1-\gamma_0)(\tilde{m}-\tilde{E})(-V_\omega^2+\bar{V}_\pi^2+2V_\omega\bar{V}_\pi) \\ -2\mathbf{k}\cdot\boldsymbol{\gamma}[V_\omega(V_\omega-V_\sigma)-3\bar{V}_\pi^2+\bar{V}_\pi(V_\sigma-4V_\omega)]\} . \tag{2.54}$$

The total density matrix is then $d(k) = d_D(k) - d_E(k)$.

Using the above equations, we can evaluate $\delta\rho(k)$. It is convenient to use the identities

$$\bar{f}_s(k)(A+\not{B})f_{s'}(k) = \delta_{ss'}\frac{E\tilde{m}}{\tilde{E}m_N}\left[A+\frac{B\cdot\tilde{k}}{\tilde{m}}\right] , \tag{2.55}$$

$$\bar{h}_s(k)(A+\not{B})h_{s'}(k) = -\delta_{ss'}\frac{E\tilde{m}}{\tilde{E}m_N}\left[A+\frac{B\cdot\tilde{k}'}{\tilde{m}}\right] , \tag{2.56}$$

$$\bar{f}_s(k)(A+\not{B})h_{s'}(k) = \langle s|\boldsymbol{\sigma}|s'\rangle\cdot\frac{E}{\tilde{E}}\left[\left[\frac{\mathbf{k}\cdot\mathbf{B}}{\tilde{E}+\tilde{m}}-A\right]\frac{\mathbf{k}}{\tilde{m}}-\frac{\tilde{E}\mathbf{B}}{\tilde{m}}\right] , \tag{2.57}$$

where $\tilde{k}' = (-\tilde{E}, \mathbf{k})$. It is also convenient, and perhaps instructive, to express the nuclear matter density matrix in the form

$$\delta\rho^{++}(k)=\sum_{s,s'} f_s(k)[D_D^{++}(k)-D_E^{++}(k)]_{ss'}\bar{f}_{s'}(k)\;, \tag{2.58}$$

$$\delta\rho^{+-}(k)=\sum_{s,s'} f_s(k)[D_D^{+-}(k)-D_E^{+-}(k)]_{ss'}\bar{h}_{s'}(k)\;, \tag{2.59}$$

$$\delta\rho^{-+}(k)=\sum_{s,s'} h_s(k)[D_D^{-+}(k)-D_E^{-+}(k)]_{ss'}\bar{f}_{s'}(k)\;, \tag{2.60}$$

$$\delta\rho^{--}(k)=\sum_{s,s'} h_s(k)[D_D^{--}(k)-D_E^{--}(k)]_{ss'}\bar{h}_{s'}(k)\;. \tag{2.61}$$

If we use this form, a direct quantitative comparison can be made with the leading term $\rho_0(k)$ of Eq. (2.6). We may obtain analytic expressions for the functions $D(k)$ in terms of the meson potentials:

$$[D_D^{++}(k)]_{ss'}=2\delta_{ss'}W(k)\frac{\tilde{m}^2}{\tilde{E}}\left[\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right]^2(V_\sigma-V_\omega)^2+3V_\omega^2+2\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right]V_\omega(V_\omega-V_\sigma)+48\left[\frac{\tilde{E}}{\tilde{m}}V_\pi\right]^2\right]\;, \tag{2.62}$$

$$[D_D^{+-}(k)]_{ss'}=2\langle s|\boldsymbol{\sigma}\cdot\mathbf{k}|s'\rangle W(k)\frac{\tilde{m}}{\tilde{E}}\left[\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right](V_\sigma-V_\omega)^2-3V_\omega^2-2\left[\frac{\tilde{m}}{\tilde{E}-\tilde{m}}\right]V_\omega(V_\omega-V_\sigma)\right]\;, \tag{2.63}$$

$$[D_D^{--}(k)]_{ss'}=2\delta_{ss'}W(k)\frac{\mathbf{k}^2}{\tilde{E}}[(V_\sigma-V_\omega)^2+3V_\omega^2-2V_\omega(V_\omega-V_\sigma)]\;. \tag{2.64}$$

Note that $D_{ss',D}^{-+}$ is the Hermitian conjugate of $D_{ss',D}^{+-}$. Since $D_{ss',D}^{+-}$ is itself Hermitian, we obtain $D_{ss'}^{-+}=D_{ss'}^{+-}$ for the direct terms. The exchange terms are more complicated, including π-ω and π-σ mixing. (Such terms do not appear in the direct amplitudes.) We have

$$[D_E^{++}(k)]_{ss'}=\delta_{ss'}W(k)\frac{\tilde{m}^2}{2\tilde{E}}\left[\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right]^2(V_\sigma-V_\omega)^2-3V_\omega^2-2\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right]V_\omega(V_\omega-V_\sigma)+48\left[\frac{\tilde{E}}{\tilde{m}}V_\pi\right]^2\right.$$
$$\left.-8\frac{\tilde{E}(\tilde{E}+\tilde{m})}{\tilde{m}(\tilde{E}-\tilde{m})}V_\pi(V_\omega-V_\sigma)-24\frac{\tilde{E}}{\tilde{m}}V_\pi V_\omega\right]\;, \tag{2.65}$$

$$[D_E^{--}(k)]_{ss'}=\frac{1}{2}\delta_{ss'}W(k)\frac{\mathbf{k}^2}{\tilde{E}}[(V_\sigma-V_\omega)^2-3V_\omega^2+2V_\omega(V_\omega-V_\sigma)]\;, \tag{2.66}$$

$$[D_E^{+-}(k)]_{ss'}=\frac{1}{2}\langle s|\boldsymbol{\sigma}\cdot\mathbf{k}|s'\rangle W(k)\frac{\tilde{m}}{\tilde{E}}\left[\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right](V_\sigma-V_\omega)^2+3V_\omega^2+2\left[\frac{\tilde{m}}{\tilde{E}-\tilde{m}}\right]V_\omega(V_\omega-V_\sigma)\right.$$
$$\left.-4\frac{\tilde{E}}{\tilde{m}}V_\pi(V_\omega-V_\sigma)+12\frac{\tilde{E}}{\tilde{m}}V_\pi V_\omega\right]\;. \tag{2.67}$$

Combining the above expressions, we can write the full density matrix for nuclear matter in a relatively simple form:

$$\rho(k)=\{D^+(k)+\Theta(k_F-|\mathbf{k}|)\delta[k^0-\epsilon^+(\mathbf{k})]\}\sum_s f_s(\mathbf{k})\bar{f}_s(\mathbf{k})$$
$$+D^0(k)\sum_{ss'}[h_s(\mathbf{k})\langle s|\boldsymbol{\sigma}\cdot\mathbf{k}|s'\rangle\bar{f}_{s'}(\mathbf{k})+f_s(\mathbf{k})\langle s|\boldsymbol{\sigma}\cdot\mathbf{k}|s'\rangle\bar{h}_{s'}(\mathbf{k})]+D^-(k)\sum_s h_s(\mathbf{k})\bar{h}_s(\mathbf{k})\;. \tag{2.68}$$

In our calculations we have obtained asymptotic values for the functions $D(k)$. These are given by the following expressions:

$$D^+(k)=W(k)\frac{\tilde{m}^2}{2\tilde{E}}\left[3\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right]^2(V_\sigma-V_\omega)^2+15V_\omega^2+10\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right]V_\omega(V_\omega-V_\sigma)+144\left[\frac{\tilde{E}}{\tilde{m}}V_\pi\right]^2\right.$$
$$\left.+8\frac{\tilde{E}(\tilde{E}+\tilde{m})}{\tilde{m}(\tilde{E}-\tilde{m})}V_\pi(V_\omega-V_\sigma)+24\frac{\tilde{E}}{\tilde{m}}V_\pi V_\omega\right]\;, \tag{2.69}$$

$$D^-(k)=\frac{1}{2}W(k)\frac{\mathbf{k}^2}{\tilde{E}}[3(V_\sigma-V_\omega)^2+15V_\omega^2-10V_\omega(V_\omega-V_\sigma)]\;, \tag{2.70}$$

$$D^0(k)+\frac{1}{2}|\mathbf{k}|W(k)\frac{\tilde{m}}{\tilde{E}}\left[3\left[\frac{\tilde{E}+\tilde{m}}{\tilde{E}-\tilde{m}}\right](V_\sigma-V_\omega)^2-15V_\omega^2-10\left[\frac{\tilde{m}}{\tilde{E}-\tilde{m}}\right]V_\omega(V_\omega-V_\sigma)\right.$$
$$\left.+4\frac{\tilde{E}}{\tilde{m}}V_\pi(V_\omega-V_\sigma)-12\frac{\tilde{E}}{\tilde{m}}V_\pi V_\omega\right]\;. \tag{2.71}$$

It should be kept in mind that although the above expressions for the $D(k)$ are defined for $|\mathbf{k}|>k_F$, they are only true asymptotically for $|\mathbf{k}|\to\infty$ and are approximately correct for $|\mathbf{k}|>2k_F$.

III. RESULTS AND DISCUSSION

The expressions for the various parts of the density matrix of nuclear matter provides a clear description of the relative importance of π, σ, and ω contributions, as well as the total magnitude of the correction to the density matrix. The importance of the contribution of the exchange diagrams should be stressed, since in some cases they include processes that vanish in the direct diagrams. Moreover, because the meson potentials are such that $V_\sigma \approx V_\omega$, there are many cancellations between the various terms, and often the result is smaller than the individual terms.

To perform numerical calculations, we use meson-nucleon form factors of the following form:

$$F_i(q^2)=\frac{\Lambda_i^2-m_i^2}{\Lambda_i^2-q^2} . \tag{3.1}$$

The meson masses are taken to be $m_\sigma=0.55$ GeV, $m_\omega=0.783$ GeV, and $m_\pi=0.138$ GeV. The parameters of the vertex form factors are $\Lambda_\sigma=1.8$ GeV, $\Lambda_\omega=1.5$ GeV, and $\Lambda_\pi=1.3$ GeV. Since $|\mathbf{p}_1|$ and $|\mathbf{p}_2|$ are small, we obtain $q^2=2\tilde{m}(\tilde{m}-\tilde{E})<0$, so that $F^{(i)}(Q^2)$ is always positive. The meson-nucleon coupling constants are taken to be $\alpha_\sigma=g_\sigma^2/4\pi=6.00$, $\alpha_\omega=g_\omega^2/4\pi=10.00$, and $\alpha_\pi=f_\pi^2(\tilde{m}/m_\pi)^2/4\pi=1.19$. The values chosen here for α_σ and α_ω differ somewhat from the values adopted in the Walecka model.[6] This is due to the fact that the Walecka values ($\alpha_\sigma=8.76$, $\alpha_\omega=11.27$) lead to an unphysically large depletion of the occupied states in the Fermi sea. We therefore treat these couplings as parameters which may be adjusted to obtain acceptable values for the density matrix. If we consider only the σ and ω mesons, for the values of the coupling constants adopted in this work, we find that the depletion of the Fermi sea is of the order of 10% (integrating for $|\mathbf{k}|\geq 0.5$ GeV). That result is in agreement with the work of Remos, Polls, and Dickoff[1] who studied depletion of the Fermi sea via a realistic (central) interaction making use of a self-consistent Green's-function approach. We may only obtain reasonable results, however, if we restrict our calculation to the σ and ω mesons. We find that the addition of pions leads to very large values for $\delta\rho(k)$ and, correspondingly, very large depletion of the Fermi sea. We conclude that, for the case of pions, the correlation effects cannot be neglected or emulated by an appropriate choice of the pion-nucleon coupling constant.

In Fig. 2 the various contributions of the mesons σ, ω, and π to D^{++} are plotted as a function of $|\mathbf{k}|$. Note that, although the contributions of the σ-σ, σ-ω, and ω-ω terms are large (compared to that of the pion or pion cross terms), they add up to a value less than 0.01. Given that cancellation, we see that dominant contribution becomes that of the pion. The importance of the pion is accentuated at large $|\mathbf{k}|$, since the pion coupling to nucleons is linear in the momentum transfer. As expected, for $|\mathbf{k}|<0.5$ GeV the approximation of neglecting $|\mathbf{p}_1|$ and $|\mathbf{p}_2|$ in the integrand cannot be trusted to give correct results, and indeed, an extrapolation of our model for $|\mathbf{k}|<0.5$ GeV results in unacceptably large values of $\delta\rho(k)$.

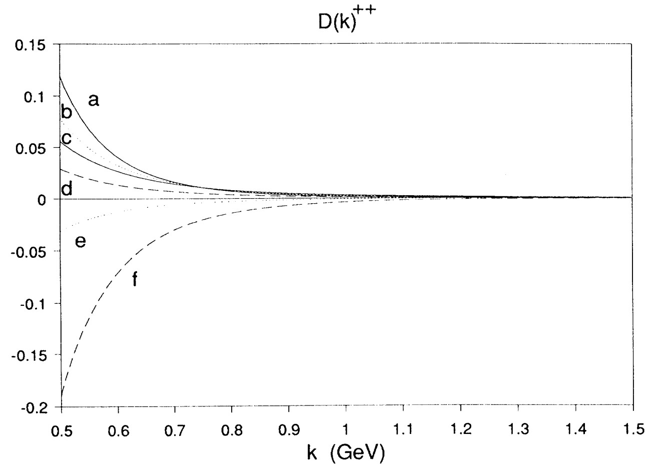

FIG. 2. Contributions to $D^{++}(k)$ from various combinations of exchanged mesons (see Fig. 1). [Note the almost complete cancellation of the contributions a, b, d, e, and f. Because of this cancellation, $D^{++}(k)$ is quite sensitive to short-range correlation effects.] (a) σ-σ, (b) ω-ω, (c) π-π, (d) π-$\omega+\omega$-π, (e) π-$\sigma+\sigma$-π, and (f) σ-$\omega+\omega$-σ.

In Fig. 3 we show the contributions of the σ, ω, and π mesons to D^{+-} as a function of $|\mathbf{k}|$. As before, the contributions of the σ and ω mesons are dominant, but they are much smaller than in the case of D^{++}. Also, because of the form of the energy denominator, these contributions fall off more slowly than in the case of D^{++}. Since D^{+-} does not contribute in the calculation of the baryon

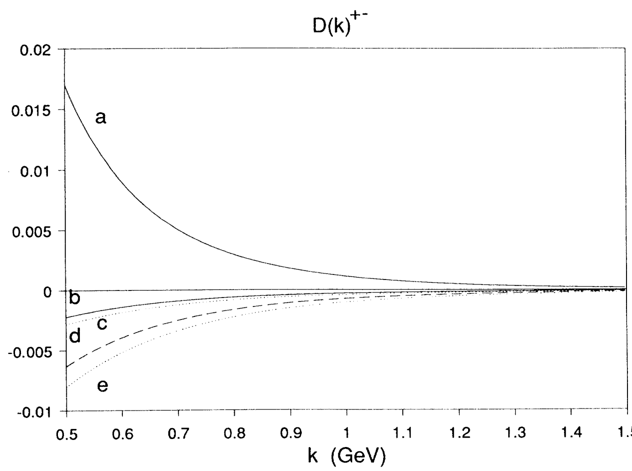

FIG. 3. Contributions to $D^{+-}(k)$ from various combinations of exchanged mesons. (a) σ-σ, (b) π-$\sigma+\sigma$-π, (c) π-$\omega+\omega$-π, (d) ω-$\sigma+\sigma$-ω, and (e) ω-ω.

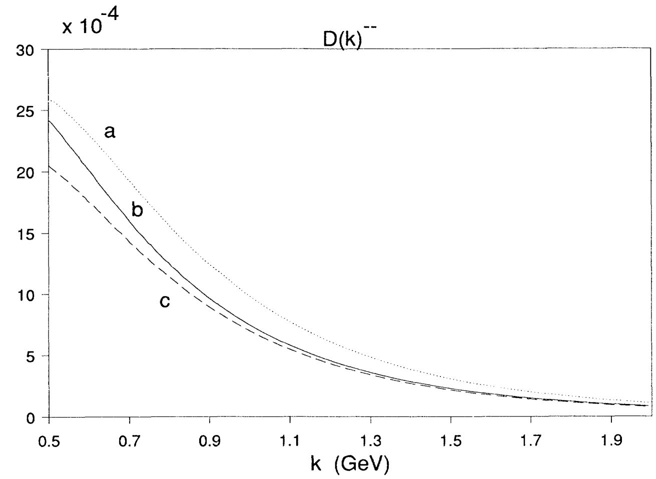

FIG. 4. Contributions to $D^{--}(k)$ from various combinations of exchanged mesons. (a) ω-ω, (b) σ-σ, and (c) ω-$\sigma+\sigma$-ω. Note that there is no contribution from pion exchange in this case and that the σ and ω exchanges are *coherent* in their contributions to $D^{--}(k)$.

density, it is not necessary that D^{+-} be small for the theory to be convergent.

The contributions to D^{--} are shown in Fig. 4. Note that only the σ and ω mesons contribute in this case [see Eqs. (2.64) and (2.67)]. Since D^{--} has an extra factor of $(\mathbf{k}/\tilde{m})^2$ relative to D^{++}, arising from the energy denominator, its contributions are small for $|\mathbf{k}|=0.5$ GeV, but dominate for $|\mathbf{k}|\geq 1.5$ GeV. Moreover, *all mesonic contributions add coherently* in D^{--}, in contrast to the situation in the calculation of D^{++}.

In Fig. 5 we show the sum of the contributions of the various mesons to D^{++}, D^{+-}, and D^{--} for the model which includes the pion contribution. In Fig. 6 we show the same curves for the model excluding pions. As can be seen from Fig. 6, D^{--} is much larger than D^{++} and thus leads to an unacceptable depletion of the occupied states in the Fermi sea. [We have shown in previous work[7] that the correlation effects are *very small,* if one calculates the low-momentum transfer amplitude for the process shown on the left-hand (or right-hand) side of the diagrams of Fig. 1. Such amplitudes may be combined to form D^{--}. However, for the range of $|\mathbf{k}|$ of interest here ($|\mathbf{k}|>0.5$ GeV), it is uncertain how important the correlation effects might be.] An alternative solution to the problem of very large values of D^{--} might be that the meson-antinucleon vertex form factors (which were taken to be equal to the meson-nucleon vertex form factors) are strongly suppressed at high-momentum transfer. This suggestion would be consistent with the suppression of antinucleon effects advocated by Brodsky.[9] Note also that antinucleons have been dropped from the analysis in Ref. 7 by systematically neglecting terms involving the negative-energy spinors $h_s(\mathbf{p})$ in the analysis. The inclusion of such terms is not necessary if one is constructing a phenomenological model, such as that described in Ref. 7. It is unclear whether any improvement is to be obtained in the relativistic description, if terms involving

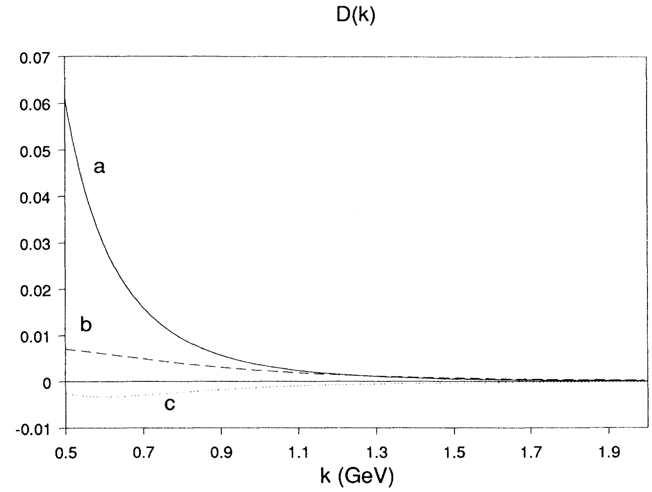

FIG. 5. Total contribution of meson exchange, including the contribution due to pions. (a) $D^{++}(k)$, (b) $D^{--}(k)$, and (c) $D^{+-}(k)$.

$h_s(\mathbf{p})$ are retained in the theory.

Finally, we note that we can introduce a quantity $B(k)$, which is proportional to the spin trace of the density matrix

$$B(|\mathbf{k}|)=\frac{N}{2}\mathrm{Tr}[\gamma_0\rho(|\mathbf{k}|)m_N/E] \tag{3.2}$$

$$=N\{D^{++}(k)+D^{--}(k)+\Theta(k_F-|\mathbf{k}|)\delta[k^0-\epsilon^+(\mathbf{k})]\} \ . \tag{3.3}$$

Note that D^{+-} and D^{-+} do not appear in Eq. (3.3), since the nondiagonal parts of $\rho(k)$ do not contribute to the trace. Here N is the normalization constant, and it is

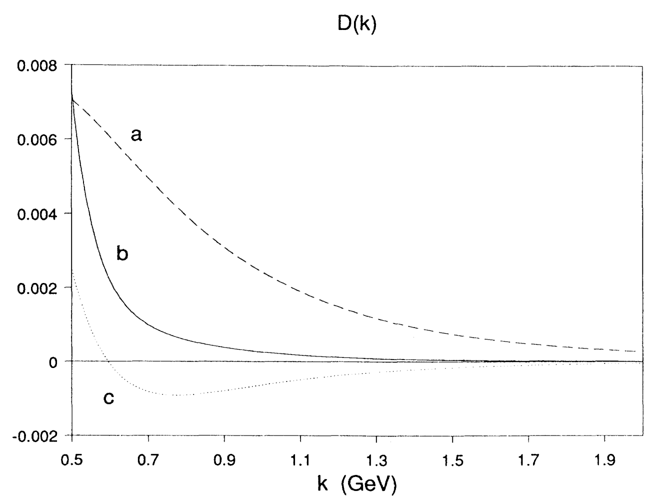

FIG. 6. Total contribution of meson exchange, excluding the contribution due to pions. (a) $D^{--}(k)$, (b) $D^{++}(k)$, and (c) $D^{+-}(k)$. Note the change of scale relative to Fig. 5. The fact that $D^{++}(k)$ is about an order of magnitude smaller in the absence of pion-exchange contributions is a consequence of the results exhibited in Fig. 2.

obtained by normalizing $B(k)$ as follows:

$$\frac{1}{(2\pi)^3}\int d^3k\, B(|\mathbf{k}|)=1 \; . \tag{3.4}$$

In the absence of interaction, we have $N=22.6$ fm^3, if $k_F=1.38$ fm^{-1}. If we consider only σ and ω mesons and drop the contribution of $D^{--}(k)$ in Eq. (3.3), we find $N=21.3$ fm^3. That result represents a depletion of the Fermi sea of about 6%. However, inclusion of $D^{--}(k)$ yields $N=13.5$ fm^3. The depletion of the Fermi sea is about 40% in this case, indicating a breakdown of the perturbative approach upon inclusion of negative-energy states in the density matrix.

We recall that we were motivated to study this problem by the rather extensive discussion of high-momentum components in nuclei associated with the analysis of y-scaling experiments.[5] In order to describe such components, one requires a relativistic description of the system. The simplest analysis may be based upon the Walecka model, for which the mean-field analysis represents the leading approximation. To go beyond the mean-field approximation, we can consider the corrections to the wave function of two-particle, two-hole character. We have studied such corrections in this work, and our results can be summarized as follows. Sigma and omega mesons may be used to calculate relativistic corrections to the density matrix of nuclear matter within the context of the Walecka model. The model breaks down, however, if one considers the complete form for $\rho(k)$, including antinucleon states. This might be an indication either that the effects of correlations must be included or that the meson-antinucleon vertex form factors are strongly suppressed at high Q^2. (An attempt to also include pion-exchange contributions results in an unacceptably large depletion of the Fermi sea, indicating that a simple exchange model, without including correlations, is inadequate in the case of pion exchange.) Resolution of both of the above deficiencies requires further study of the correlation corrections to the exchange of σ, ω, π, and ρ mesons. It is unclear whether inclusion of correlation effects can suppress the quite large coupling to negative-energy states exhibited in our analysis or whether negative-energy states will have to be suppressed by including new form factors in the analysis.[9] However, our past experience leads us to believe that short-range correlation effects will be quite unimportant in suppressing the excitation of negative-energy states.[7] This observation may be understood by noting that correlation effects are important for D^{++} because of an almost complete cancellation between σ and ω exchange. Therefore, small corrections arising from correlations can make large changes in D^{++}. On the other hand, σ and ω exchanges are coherent in the calculation of D^{--}, and therefore correlations will only modify D^{--} to a small extent. We hope to continue our studies of these issues in a future work.

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