presumably enhance $M1$ transitions. It seems, therefore, that the signature of nonstatic spin-isospin order is the lowering of the energy and the enhancement of the electromagnetic transition amplitudes of the levels of unnatural parity. This is considered a precursor to pion condensation,[8] i.e., to static spin-isospin order. Precursor phenomena are usually studied in the framework of the random-phase approximation, whose close relation to our semiclassical description has already been emphasized. The exact connection between precursor phenomena and the present model, however, is not straightforward and will be investigated separately. We further observe that correlations other than spin-isospin will occur in the ground state. Coexistence of different zero-point modes will, of course, reduce the transition amplitudes relative to our estimates.

Let us finally note that $K_{\rm OPE}$ is strongly increasing with the density. For instance, $K_{\rm OPE}$ increases from -13 to -19 MeV fm^{-2}, increasing the density by 20% in ^{12}C. Spin-isospin oscillations can become more pronounced with increasing density, and therefore acquire relevance to heavy-ion collisions.

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Energy Dependence of the Coupling Potentials in (p,n) Reactions

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We investigate the physical origin of the energy dependence of the isovector potentials V_τ and $V_{\tau\sigma}$ within a model which includes the one-pion and one-ρ—exchange potential as well as the second-order effects of the tensor force.

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In the past year, new experimental information has been obtained on collective spin resonances in nuclei using highly energetic protons in (p,n) reactions.[1,2] The physical origin of this new development is connected with a strong energy dependence of the effective V_τ-coupling potential which has been found in those experiments. At incident proton energies below 50 MeV the "electric" charge-exchange resonances, like the well-known isobaric-analog resonances (IAR), which are connected with the isospin operator $\vec{\tau}$ only, dominate the experimental cross sections. The "magnetic" charge-exchange resonances, which are connected with the spin and isospin exchange operators $\vec{\sigma}$ and $\vec{\tau}$, are much less pronounced.[3] The situation is reversed if the energy of the incoming proton is increased. In the case of 160-MeV protons the structure of the differential cross section at 0° of the reaction $^{90}Zr(p,n)^{90}Nb$ is qualitatively different compared to the low-energy data. It shows a very strong new resonance which has been interpreted as the Gamow-Teller (GT) resonance (1^+) because it possesses a $\Delta l = 0$ angular distribution.[2] The IAR shows up at this high proton energy only as a little peak at the lower-energy side of the GT resonance. All the data show that the strength of the analog transitions relative to the new resonances strongly decreases with increasing proton energy. From a preliminary analysis of the experimental spectra with use of the distorted-wave Born approximation (DWBA), it follows that the ratio of the corresponding coupling potentials $[V_{\sigma\tau}(0)/V_\tau(0)]^2$ increases by a factor of about 10 if one changes the proton energy from 40 to 200 MeV.[4] Austin *et al.*[5] investigated the reaction $^7Li(p,n)^7Be$ using 25- and 45-MeV protons. Here the ratio $[V_{\sigma\tau}(0)/V_\tau(0)]^2$ increases by about 60%. Petrovich and Love obtained this unexpected energy dependence of the coupling potentials from their G-matrix interaction[6] in good qualitative agreement with the experimental findings. In this note we investigate the physical origin of the energy dependence of these isovector coupling potentials.

In the following we are interested in the $T = 1$ channel of the effective nucleon-nucleon interaction. In addition, we restrict ourselves to forward scattering ($\vec{k}_1 = \vec{k}_3$ in Fig. 1) where the contribution of the tensor force is small. Therefore we will calculate only the central parts of the coupling potential with the additional condition $\vec{k}_1 = \vec{k}_3$ (scattering angles $\theta = 0$):

$$V(\theta = 0) = V(0) = V_\tau \vec{\tau}_1 \cdot \vec{\tau}_2 + V_{\sigma\tau} \vec{\tau}_1 \cdot \vec{\tau}_2 \vec{\sigma}_1 \cdot \vec{\sigma}_2. \quad (1)$$

We start with a model of the isovector force which has been used extensively in connection with pion physics[7-9] and magnetic properties of nuclei.[10] The interaction has the following simple structure: (i) We consider explicitly the one-pion exchange and one-ρ exchange, because these mesons carry the isospin one. (ii) The effects of the other mesons (especially of the ω meson) are summarized in a pair correlation function.

For the isospin-1 channel, our interaction has the same content as the Paris potential,[11] aside from short-range terms in the latter arising from A_1 exchange and small terms remaining from the nucleon box diagrams, once the iterated one-

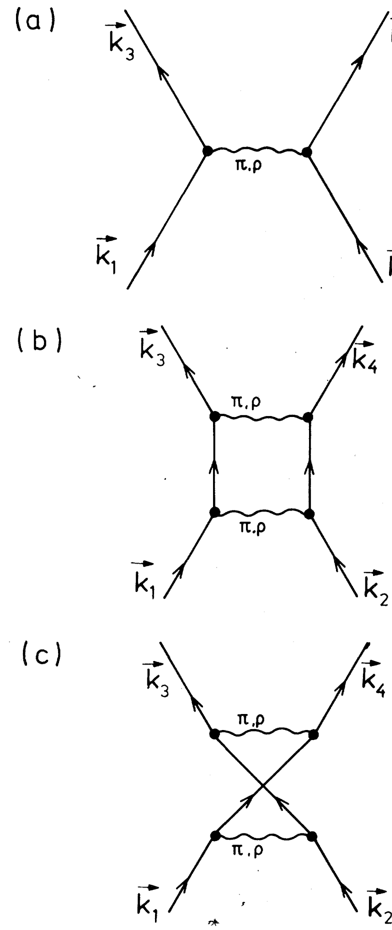

FIG. 1. First- and second-order contributions of the one-pion and one-ρ exchange included in the present model.

pion–exchange potential (OPEP) has been removed. In particular, we shall use a strong ρ-meson tensor coupling, corresponding to the same $J = 1$ helicity amplitudes[12] as in Ref. 11. Of course, our pion coupling is the same. A crucial point is that the large box diagrams involving intermediate isobars have only small isovector parts.[13] In this respect, the isovector interaction is much simpler than the isoscalar one, and the calculation of the effective interaction from the elementary particle exchange can be made rather direct, as we shall show.

We begin, then, with the following form:

$$W = W(q)\vec{\sigma}_1 \cdot \vec{\sigma}_2 \vec{\tau}_1 \cdot \vec{\tau}_2 = \int [d^3k/(2\pi)^3] g(\vec{q} - \vec{k})[V_\pi(\vec{k}) + V_\rho(\vec{k})], \quad (2)$$

where V_π and V_ρ are the bare one-pion– and one-

ρ-exchange potentials, respectively. If we replace the exact correlation function $g(\vec{q}-\vec{k})$ by its dominant Fourier component,[8]

$$g(\vec{q}-\vec{k})=(2\pi)^3\delta(\vec{q}-\vec{k})-(2\pi^2/q_c^2)\delta(|\vec{q}-\vec{k}|-q_c), \tag{2a}$$

with $q_c=3.93\ \mathrm{fm}^{-1}$, which is of the order of the ω-meson mass, we obtain, e.g., for the central part of V_π,

$$W_{\mathrm{dir}}^{\pi}(q)=\frac{4\pi f_\pi^2}{m_\pi^2}\vec{\tau}_1\cdot\vec{\tau}_2\vec{\sigma}_1\cdot\vec{\sigma}_2\frac{1}{3}\left\{\frac{(\Lambda_\pi^2-m_\pi^2)^2}{(\Lambda_\pi^2+q_c^2+q^2)^2-4q_c^2q^2}+\frac{m_\pi^2(\Lambda_\pi^2-m_\pi^2)}{(\Lambda_\pi^2+q_c^2+q^2)^2-4q_c^2q^2}\right.$$
$$+\frac{m_\pi^2}{4q_cq}\ln\left[\frac{(m_\pi^2+q_c^2+q^2-2q_cq)(\Lambda_\pi^2+q_c^2+q^2+2q_cq)}{(\Lambda_\pi^2+q_c^2+q^2-2q_cq)(m_\pi^2+q_c^2+q^2+2q_cq)}\right]$$
$$\left.-\frac{q^2}{q^2+m_\pi^2}\left(\frac{\Lambda_\pi^2-m_\pi^2}{\Lambda_\pi^2+q^2}\right)^2\right\}, \tag{2b}$$

$$W_{\mathrm{dir}}(q)=W_{\mathrm{dir}}^{\pi}(q)+W_{\mathrm{dir}}^{\rho}(q), \tag{2c}$$

where we obtain W_{dir}^{ρ} from Eq. (2b) by replacing m_π and Λ_π by m_ρ and Λ_ρ, respectively. The cutoff parameters were chosen to be $\Lambda_\pi=1.2$ GeV and $\Lambda_\rho=2$ GeV, $f_\pi^2=0.081$ and $f_\rho^2=4.86$ (strong ρ-meson coupling). It is obvious that W_{dir} gives only a contribution in the $\sigma\tau$ channel and that for $\vec{q}=\vec{k}_3-\vec{k}_1=0$ (see Fig. 1) this term is energy independent (i.e., independent of k_1). The (isovector) exchange terms to Eq. (2) follow immediately by applying the exchange operator $P_{\mathrm{ex}}=-P_\sigma P_\tau P(\vec{k}_3\rightarrow\vec{k}_4)$. If we use, in addition, the forward scattering condition $\vec{k}_1=\vec{k}_3$ and $\vec{k}_2=\vec{k}_4$, we get

$$W_{\mathrm{exch}}=W_{\mathrm{dir}}(|\vec{k}_2-\vec{k}_1|)(3\vec{\tau}_1\cdot\vec{\tau}_2-\vec{\sigma}_1\cdot\vec{\sigma}_2\vec{\tau}_1\cdot\vec{\tau}_2)\cdot\tfrac{1}{4}. \tag{3}$$

This part of the force is first of all energy dependent (it decreases with increasing $|\vec{k}_2-\vec{k}_1|$) and it gives a contribution in the τ channel, the magnitude of which, however, is too small. On the other hand, we know that second-order effects of the tensor force are large.[14] Therefore we also include in our effective coupling potential the effects of the second-order tensor contribution which is shown in Figs. 1(b) and 1(c). The isovector part of the direct term is of the form

$$U_{T,\mathrm{dir}}^{(2)}=-(12\vec{\tau}_1\cdot\vec{\tau}_2+4\vec{\tau}_1\cdot\vec{\tau}_2\vec{\sigma}_1\cdot\vec{\sigma}_2)U_T(q) \tag{4}$$

with

$$U_T(q)=4\mathrm{m}\int_0^\infty dz\,\frac{z^2f(z)}{q^2-4z^2}, \tag{4a}$$

where $q=|\vec{k}_2-\vec{k}_1|$ and

$$f(z)=2\pi\int_{-1}^{+1}dx\,|V^T[\tfrac{1}{2}(q^2+4z^2+4zq\cdot x)^{1/2}]|^2. \tag{4b}$$

In calculating these second-order processes, we employ the following simple tensor interaction $V^T(x)$. The long-range part is given by OPEP. The ρ-exchange tensor interaction has the same functional form as OPEP, but with m_π replaced everywhere by m_ρ, and f_π^2 by $-f_\rho^2$.[15] The fact that the ρ-exchange interaction has opposite sign to OPEP means that it will tend to cut the latter off with decreasing r; in fact, with our f_ρ^2 the combined interaction goes through zero at $r\simeq0.6$ fm. Inside this distance, the interaction is very uncertain. (In the Paris potential, it is replaced here by a weak phenomenological term.) Repulsion from ω-meson exchange is large in this inner region, making the wave function small here. Therefore, we set the tensor interaction zero in this inside region. We can then obtain the Fourier transform analytically.[16] The real part of the expression $U_T(q)$ depends strongly on the energy of the incoming particle $\sim k_1^2$ because Eq. (4a) is of the form of a principal-value integral and $f(z)$ is only slightly energy dependent. In Fig. 2 the functional dependence of $z^2f(z)$ is given for two different values of k_1. The exchange term [Fig. 1(c)] is slightly more complicated but it shows qualitatively the same behavior. Our effective coupling potential is given finally as

$$V=W_{\mathrm{dir}}+W_{\mathrm{exch}}+U_{\mathrm{dir}}+U_{\mathrm{exch}}. \tag{5}$$

The final result of our calculation is given in Fig. 3 and Table I. In Fig. 3 we show $V_\tau(\theta=0)$ and $V_{\tau\sigma}(\theta=0)$ as a function of the energy E of the incoming proton. In agreement with the experimental findings the coupling potential in the $\sigma\tau$ channel is nearly energy independent whereas the corresponding quantity in the τ channel is strongly reduced with increasing proton energy. The corresponding theoretical ratio for 25- and 45-MeV protons is of the order of 40%, which is also in fair agreement with the experimental findings[5] (~60%). In Table I, we give the various

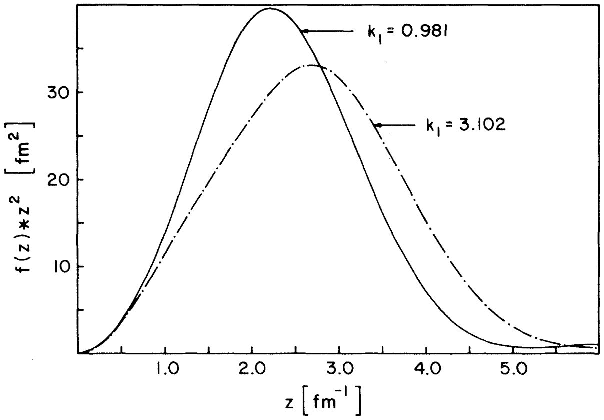

FIG. 2. Functional behavior of $z^2f(z)$ [Eq. (4b)] calculated for two different proton energies $k_1 \approx 20$ MeV and $k_1 \approx 200$ MeV.

contributions to V_τ and $V_{\tau\sigma}$ for the proton energies $E = 20$ and 130 MeV. At the lower energy about $\frac{3}{4}$ of V_τ comes from the second-order effect of the tensor and $\frac{1}{4}$ from the exchange terms of the one-pion– and one-ρ–exchange potential. The one-pion– and one-ρ–exchange potential (including the short-range correlations), on the other hand, give rise to 60% of the $V_{\tau\sigma}$ potential whereas 40% are due to the second-order tensor contribution. We note that there is no double counting between the second-order tensor contribution and the ρ meson as long as one considers the ρ meson as an elementary particle.[9] If necessary, what little double counting there is could be eliminated by the use of helicity amplitudes as in Ref. 9.

With increasing energy (k_1) not only does the

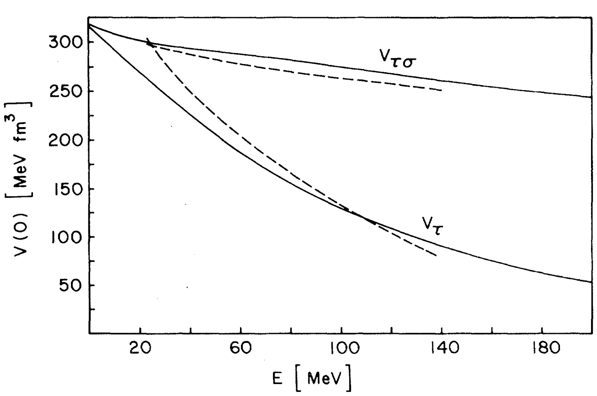

FIG. 3. Dependence of the isovector coupling potentials $V_\tau(q=0)$ and $V_{\tau\sigma}(q=0)$ on the energy of the incoming particle energy. Dashed lines give the Love-Petrovich (Ref. 6) interactions.

TABLE I. The contributions U_τ and $U_{\tau\sigma}$ (second-order tensor) and W_τ and $W_{\tau\sigma}$ (one-pion– and one-ρ–exchange potential) to the isovector coupling potentials V_τ and $V_{\tau\sigma}$ in MeV$\cdot$fm^3. The calculation has been performed for two different proton energies (20 and 130 MeV).

Proton energy	V_τ	$V_{\tau\sigma}$	U_τ	$U_{\tau\sigma}$	W_τ	$W_{\tau\sigma}$
20	270	300	210	115	60	185
130	100	265	90	60	10	205

second-order tensor contribution get smaller because of the principal-value integral, but also the Pauli exchange terms [Eq. (3)] decrease strongly. The latter one is also responsible for the increase of $W_{\tau\sigma}$.

Our results are very close to the G-matrix interactions of Petrovich and Love[6] which are shown as dashed lines in Fig. 3. The present calculation gives the physical explanation of the G-matrix result. We found within our model that the $V_{\tau\sigma}(0)$ coupling potential is nearly energy independent because it is essentially given by the (direct term) of the one-pion– and one-ρ–exchange potential which is energy independent because it is a first-order process. The only first-order direct term which contributes to the V_τ potential comes from the vector-coupling ρ-meson exchange; this is small, compared with the coupling of the ω meson, which gives a potential of the same range. In the present model we include, in addition to the Pauli terms of the one-pion– and one-ρ–exchange potentials, the second-order effects of the corresponding tensor terms which are known from other considerations to be quite large. Both contributions are strongly reduced with increasing proton energy for simple and well-known reasons.

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Electron Antineutrino Spectrum for $^{235}U(n,f)$

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The $\bar{\nu}_e$ spectrum has been computed for fission-product decay following a 30-d irradiation of ^{235}U by thermal neutrons. Estimated uncertainties lie in the range (7–15)% for $E_{\bar{\nu}} < 6$ MeV. Analysis relied on comparisons of calculated β-ray data with recently obtained experimental β-ray spectra. The $\bar{\nu}_e$ spectrum is softer than all other calculated $\bar{\nu}_e$ spectra. Cross sections (10^{-44} cm^2/fission) were calculated as $\sigma(\bar{\nu}_e + p \to n + e^+) = 58 \pm 3$, $\sigma(\bar{\nu}_e + d \to n + p + \bar{\nu}_e) = 2.7 \pm 0.2$, and $\sigma(\bar{\nu}_e + d \to n + n + e^+) = 1.04 \pm 0.13$.

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In view of ongoing neutrino oscillation experiments with $\bar{\nu}_e$ fluxes from reactors, it is important to investigate the $\bar{\nu}_e$ spectrum, and, in particular, to establish the uncertainties associated with this spectrum. The $\bar{\nu}_e$ spectrum is obtained by "summation" methods. For each and every fission product a β-ray spectrum and an associated $\bar{\nu}_e$ spectrum are calculated. All of the spectra are weighted by the amount of the responsible fission product and then summed. The calculation requires as input data (a) fission-product yields and (b) nuclear data for β decay. Particularly for the short-lived fission products, hard experimental data are incomplete or nonexistent even for the extensively studied $^{235}U(n,f)$ fission system. Consequently, much of the "input data" must be obtained from estimates, usually assumptions based upon extrapolation from known radionuclide decay, and these assumptions are different for the two previously published calculations.[1,2] In these calculations comparisons were also made with β-ray spectra measured by Tsoulfanides, Wehring, and Wyman,[3] and the calculated[1,2] β-ray data disagree with the experimental data, by as much as a factor of 2 in the high-energy portion of the spectrum. In addition, high-precision experimental β-ray spectra recently obtained by Dickens and co-workers[4-6] substantially disagree with the earlier experimental data[3] for times <10 sec following fission, particularly for $E_\beta > 4$ MeV. These experimental β-ray data,[4-6] which were obtained for short times (2.2–13 950 sec) after fission of ^{235}U, were summed to provide a spectrum for the time conditions equivalent to a pulse of fissions, a 2.2 sec cooling time, and $t_{count} = 10\,798$ sec, and this β-ray spectrum is shown in Fig. 1. The data shown provide 82% of the total energy release expected for the conditions existing after 3 h of uniform thermal-neutron fission of ^{235}U. The calculational methods of Avignone and Greenwood[2] preclude comparison with these data, and it is not evident from the text of Davis *et al.*[1] whether their calculational methods can be so tested. In any event, because of the disagreements between the two calculations and between calculations and disparate experimental data, the