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Universal Theory of Primary Interactions, Scalar Mesons, and Nucleon-Nucleon Scattering*

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The nucleon-nucleon phase shifts are calculated by adding the exchange of the scalar-isoscaler meson ϵ to our previous calculations based on the universal theory of primary interactions proposed by Sudarshan. Some improvements are obtained. In particular, the $1S_0$, $1D_2$, and $3P_2$ results are good for $g_{\epsilon NN}^2/4\pi = 13.5$ and a hard-core cutoff at 0.4 F.

I. INTRODUCTION

In a recent calculation[1] (henceforth referred as I), the universal theory of primary interactions proposed by Sudarshan[2] has been used to calculate low-energy nucleon-nucleon scattering phase shifts. The two-nucleon potential is generated by the exchange of π, η, ρ, ω, ϕ, A_1, D, and E mesons. Among these eight mesons there are two pseudoscalar, three vector, and three axial-vector mesons. In the universal theory of primary interactions, the couplings of these eight mesons are given in terms of one over-all coupling constant, which, in turn, can be determined from pion-nucleon scattering lengths and the width of the ρ decay into two pions.[3] Thus, one obtains a nucleon-nucleon potential depending on no arbitrary parameters except the masses of the exchanged mesons, if one desires to consider the masses as such. We rather take the point of view of considering them as experimental numbers.

In previous calculations,[4] it has been found that besides the exchange of a pion, which is responsible for the longest-range force, one needs to exchange vector and scalar mesons to explain the important features of nucleon-nucleon scattering. The vector-meson exchanges give rise to spin-orbit and spin-spin forces. The exchange of a scalar meson is required to provide intermediate-range attraction necessary for nuclear binding as well as phase-shift fits. No scalar meson can be accommodated in the universal theory of primary interactions without some modification. So our calculation in Paper I, which was motivated as a test of Sudarshan's theory, did not contain the exchange of scalar mesons. It had, instead, the added feature of axial-vector-meson exchanges. The exchange of a $T=0$ scalar meson will give rise to a large attractive central potential U_c in all states. There will also be a large $\mathbf{L}\cdot\mathbf{S}$ potential U_{LS} which is attractive for triplet states with $J=l+1$ and repulsive for other triplet states. The potentials U_T and $U_{\sigma p}$ will be proportional to k^2 and so will be negligible in the nonrelativistic limit. The remaining potential, U_σ, will be quite small unless the mass of the scalar meson is much higher than the nucleon mass. Now one might raise the question, can the axial-vector exchange simulate the effects of a $T=0$ scalar-meson exchange? To answer this, we examine Table I in Paper I. The D meson, which has $T=0$, has a large U_T and $U_{\mathbf{L}\cdot\mathbf{S}}$, while the other potentials are quite small. Both of these potentials have the same signs as those arising from $T=0$ scalar-meson exchange. The potentials U_σ and $U_{L\cdot\mathbf{S}}$ obtained from A_1 exchange are large and have the same (opposite) signs for $T=1$ ($T=0$) states as the $T=0$ scalar potentials. The remaining A_1 potentials are quite small. Thus for $T=1$ states the potentials arising from A_1- and E-meson exchanges add up to represent the exchange of a scalar meson. The central potential U_c will have the same sign as that of scalar-meson exchange, but with a smaller magnitude. In certain states, the largeness of U_σ for A_1 exchange will spoil the comparison with scalar-meson exchange. As, for example, in $1S_0$ and $1D_2$ states (in general singlet $T=1$ states), the net potential due to axial-vector-meson exchanges is repulsive. This is entirely due to the dominance of U_σ

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and A_1 exchange. The outcome of this is the hard core in $1S_1$ and $1D_2$, which is a desirable feature. But one has to pay a price for this by having a shallow intermediate-range attractive well, the consequence of which is a rather small slope for $1S_0$ and a too small $1D_2$ phase shift. In $T=0$ states, the dominant potentials due to A_1 exchange have signs opposite to those of a $T=0$ scalar-meson exchange. It is quite obvious that the state of affairs can be improved to some extent by including a $T=0$ scalar meson. One can, of course, include a $T=1$ scalar meson as well. The only way a scalar meson can be added to the calculation of Paper I is by hand. This will certainly introduce arbitrary parameters. We look at the problem qualitatively. A perfect fit with experiment using minimization techniques is not aimed at. To keep things simple we add the exchange of a $T=0$ scalar meson, commonly known as the ϵ meson, with a mass of 700 MeV.[5] The only arbitrary parameter, besides the cutoff, will be the coupling constant of ϵ with the nucleons.

In Sec. II we write down the expressions for the potentials due to a scalar-meson exchange. All detailed formulas may be found in Paper I. In Sec. III the results are discussed.

II. POTENTIALS OBTAINED FROM $T=0$ SCALAR-MESON EXCHANGE

The coupling of the scalar meson ϵ to nucleons is taken as

$$L=g_{\epsilon NN}\bar{\psi}\psi\phi\,, \tag{1}$$

where ψ and ϕ are the nucleon and the scalar-meson fields, respectively. The potentials, using the notations of I, are

$$U_c=-\frac{g_{\epsilon NN}^2}{4\pi}Z_0(m_\epsilon r)\left[\left(1-\frac{8m_\epsilon^2}{8m^2}\right)^2+\frac{m_\epsilon^2k^2}{16m^4}-\frac{m_\epsilon^4k^2}{128m^6}\right], \tag{2}$$

$$U_\sigma=-\frac{g_{\epsilon NN}^2}{4\pi}Z_0(m_\epsilon r)\left(\frac{m_\epsilon^4}{64m^4}+\frac{m_\epsilon^2k^2}{24m^4}-\frac{m_\epsilon^4k^2}{128m^6}\right), \tag{3}$$

$$U_T=\frac{g_{\epsilon NN}^2}{4\pi}Z_2(m_\epsilon r)\frac{m_\epsilon^2k^2}{48m^4}, \tag{4}$$

$$U_{LS}=-\frac{g_{\epsilon NN}^2}{4\pi}Z_1(m_\epsilon r)\left(\frac{m_\epsilon^2}{2m^2}-\frac{m_\epsilon^4}{16m^4}-\frac{k^2}{8m^2}+\frac{m_\epsilon^2k^2}{32m^6}\right), \tag{5}$$

and

$$U_{\sigma p}=\frac{g_{\epsilon NN}^2}{4\pi}Z_0(m_\epsilon r)\left(\frac{1}{64m^4}-\frac{k^2}{128m^6}\right), \tag{6}$$

where m_ϵ and m are the ϵ-meson and nucleon masses, respectively. These potentials are added to the potentials of Paper I with $g_{\epsilon NN}^2/4\pi$ as an adjustable parameter. Calculations are now made as in Paper I, and the results obtained are discussed in the next section. For all details of calculation techniques, the reader is referred to Paper I.

III. DISCUSSION OF RESULTS

We have calculated phase shifts for both hard-core and soft-core types of cutoff. The coupling constant $g^2/4\pi$ of the universal theory was taken as 6.5, and for each cutoff the coupling constant $g_{\epsilon NN}^2/4\pi$ was varied to obtain the best possible results. For certain states such as $3P_0$, $3D_1$, and $3F_2$, the soft- and hard-core results are quite similar. In fact, for partial waves higher than the D waves, the results should not and do not depend on the type of cutoff. The soft-core results for $3S_1$, ϵ_1, $3P_1$, $3P_2$, and ϵ_2 are bad, while for $1P_1$ and $1D_2$, the soft-core results are only slightly worse than those for a hard core. In contrast to our calculations in Paper I, which favored soft-core cutoffs, here we definitely need a hard core. The reason is quite obvious: The scalar meson introduces strong attractions that tend to cancel the hard cores present in the potentials of Paper I. We discuss below only the hard-core results. They are displayed in Fig. 1.

The introduction of the scalar meson deepens the attractive well at the intermediate range for the $1S_0$ and $1D_2$ states (in fact, the hard core given by the axial-vector meson exchanges is removed), and, as a result, the slope of the $1S_0$ phase shift increases. There seems to be, in this respect, a linear relationship between the cutoff position and the corresponding value of the coupling constant $g_{\epsilon NN}^2/4\pi$ necessary to fit the data. We get very good agreement with the experimental values of the $1S_0$ phase shifts when the cutoff is at 0.4 and $g_{\epsilon NN}^2/4\pi\approx16$. The $1S_0$ and $1D_2$ phase shifts, then, show definite improvements over I. For the unmixed P waves, on the other hand, the introduction of a scalar meson does not make things better. The $1P_1$ results seem to be worse than those in Paper I, and the $3P_0$ and $3P_1$ results are practically the same, perhaps even a little worse than in I. Some improvement is found for the $3D_2$ phase shifts.

The mixed states $3S_1$, ϵ, and $3D_1$ show slight changes from Paper I. A cutoff at 0.4 F seems to be preferred. The result for $3P_2$ with a cutoff at 0.4 F shows a definite improvement over Paper I, while ϵ_2 and $3F_2$ show very little change. The $3D_3$ results is worse, but ϵ_3 and $3G_3$ show practically no change. The results for $3F_4$ also improve, as expected. Note that both $3P_2$ and $3F_4$ are bad in Paper I for the same reason, and the introduction of a scalar meson removes the cause and both the results improve.

We find, therefore, that by adding a $T=0$ scalar-meson exchange, the $1S_0$, $1D_2$, and $3P_2$ phase shifts

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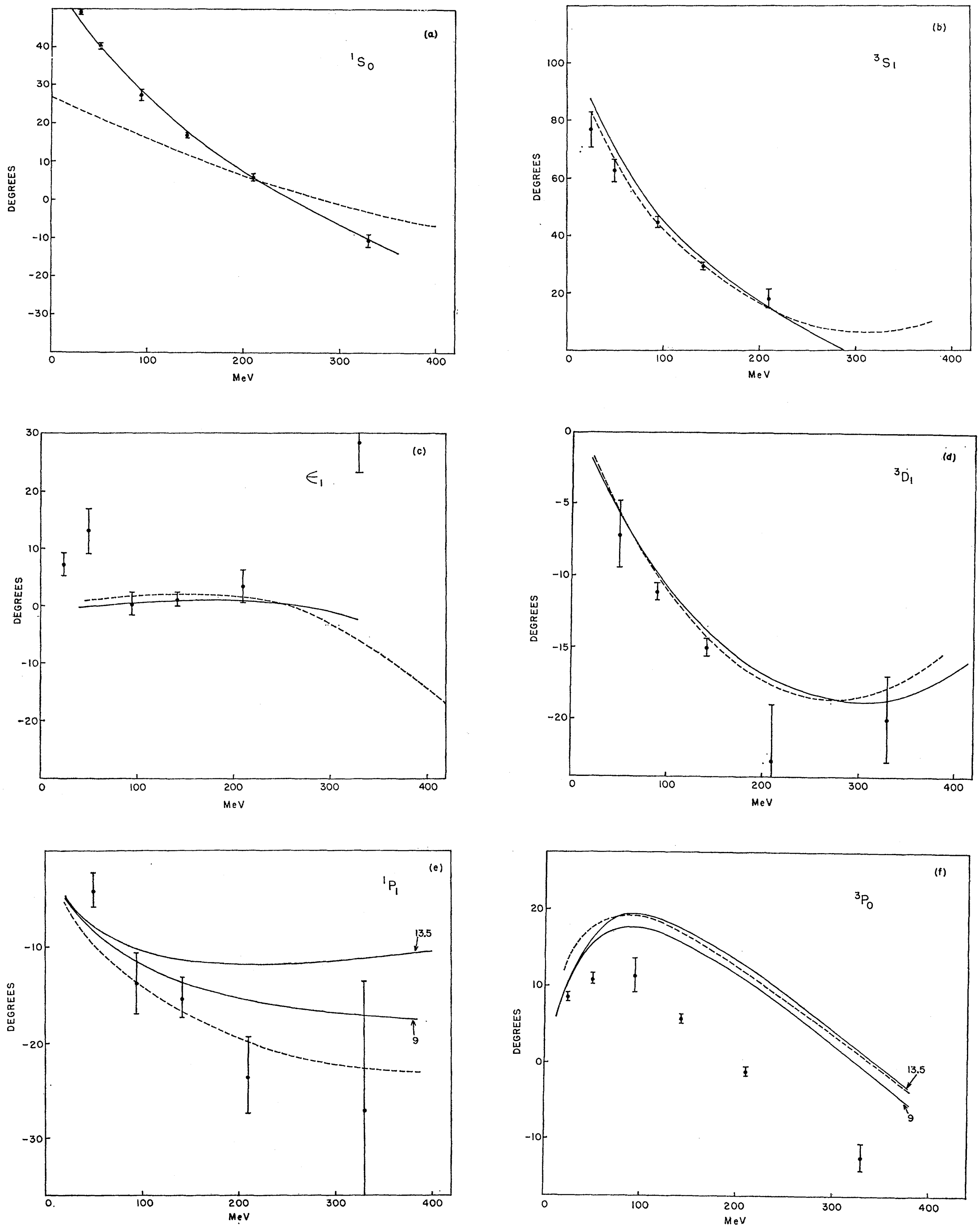

FIG. 1. Phase shifts and mixing parameters obtained in the present calculation (solid lines) plotted as a function of energy, and compared with those computed in I for a soft-core cutoff equal to 6.0 F and for a value of $g^2/4\pi$ equal to 6.5 (dotted lines). We have also indicated next to the solid lines the corresponding value of $g_{\epsilon NN}^2/4\pi$ when it differs from 13.5. A single solid line indicates no significant difference between old and new results.

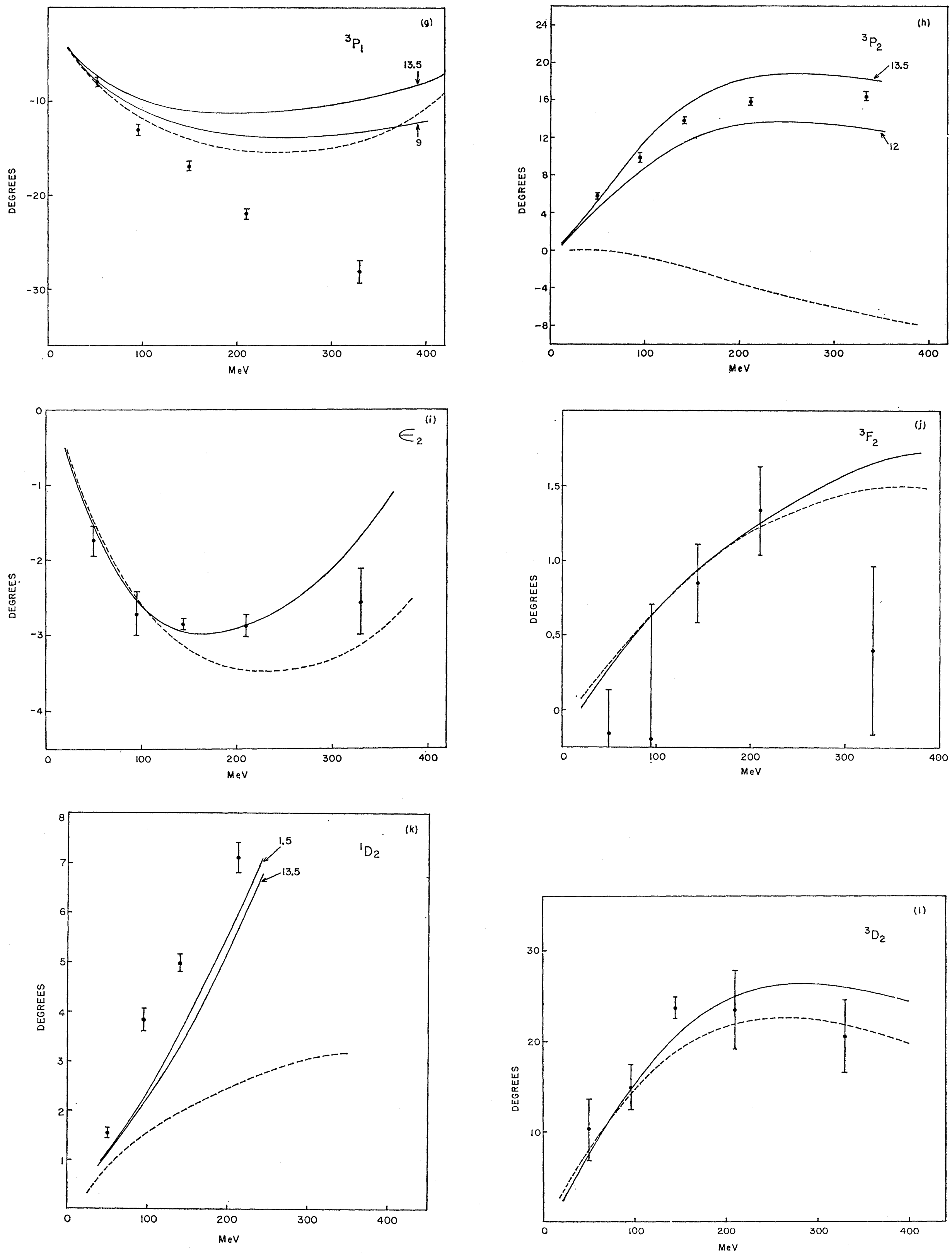

FIG. 1 (*continued*)

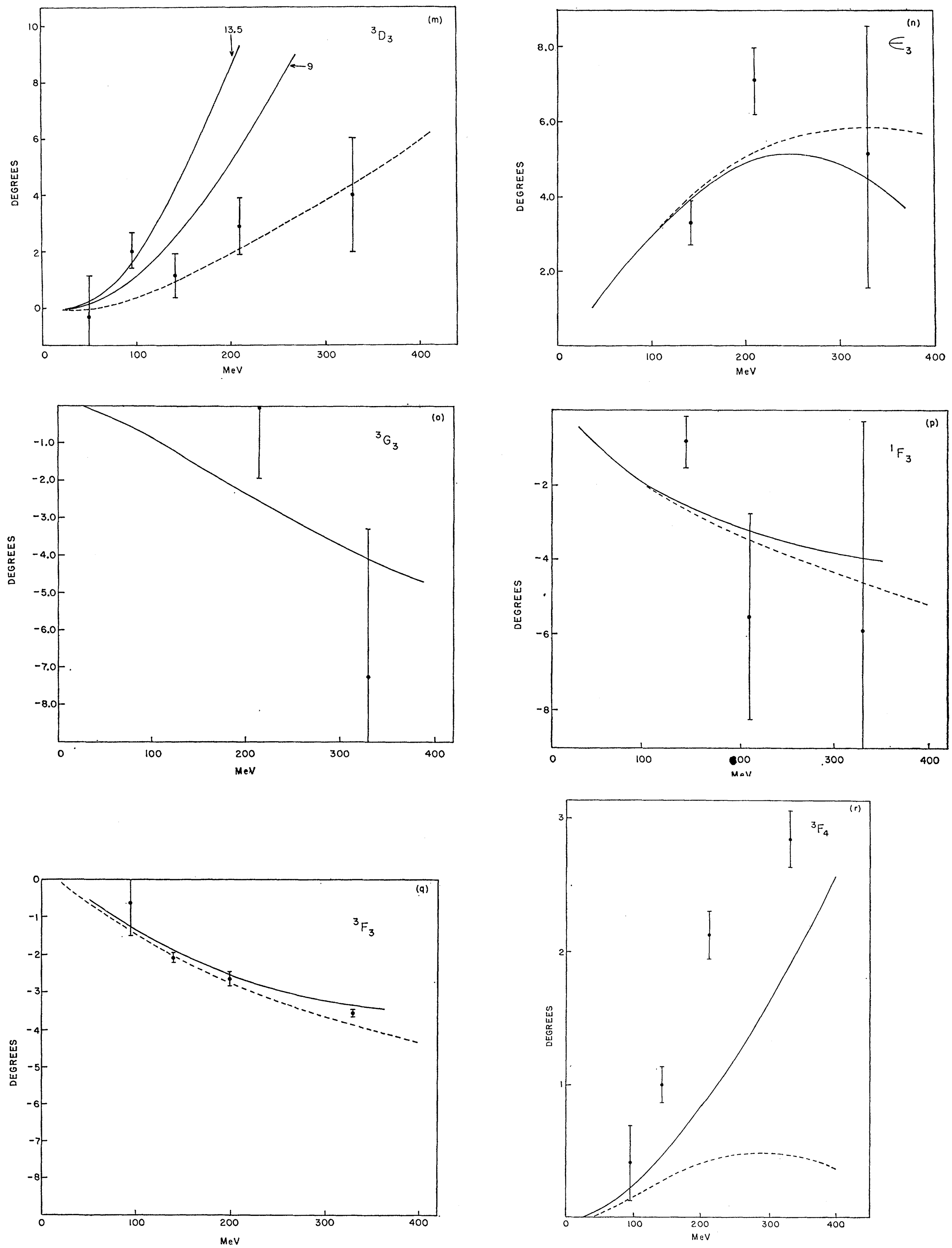

FIG. 1 (*continued*)

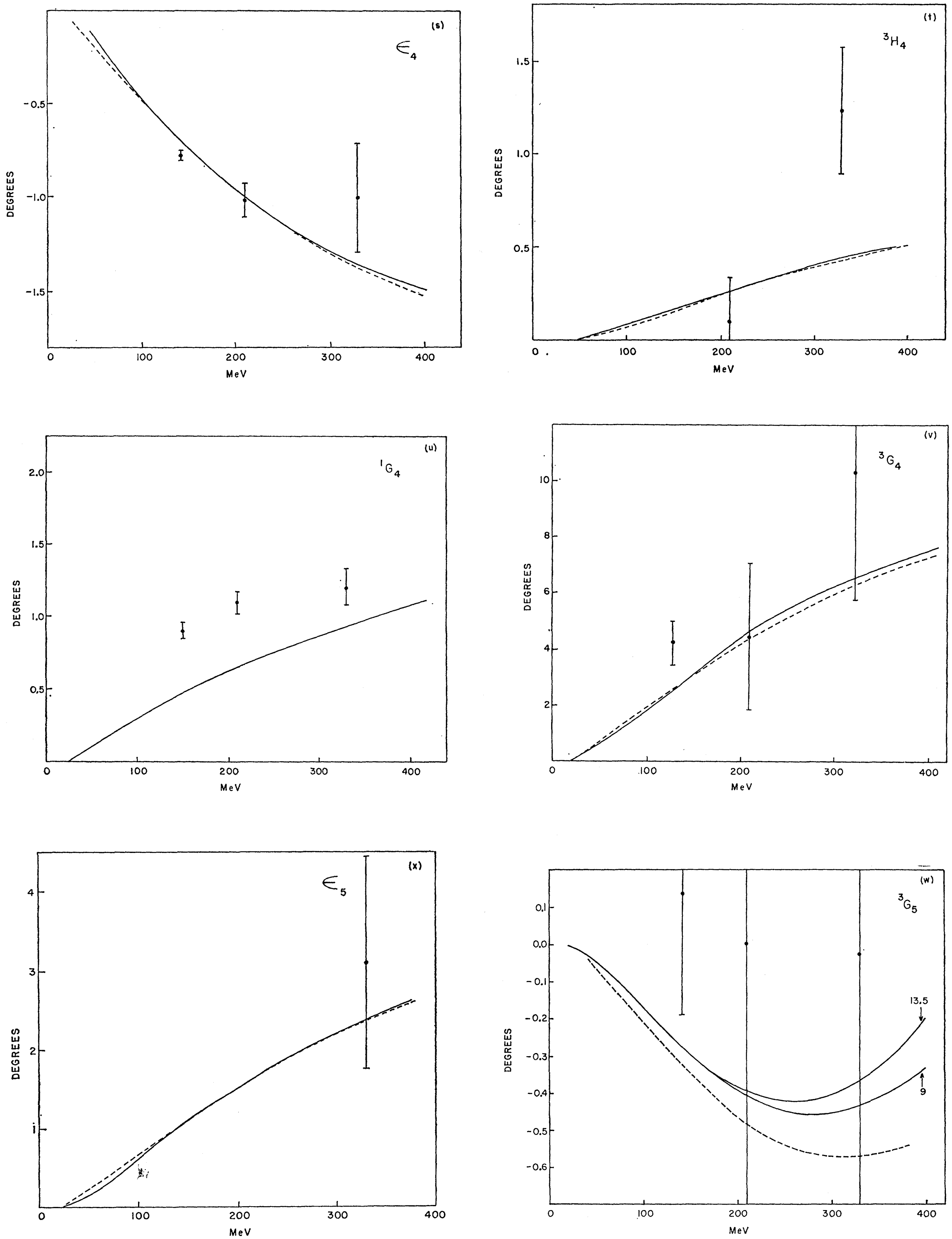

FIG. 1 (*continued*)

show a great deal of improvement over Paper I, while the P waves, with the exception of $3P_2$, get worse but not alarmingly so. The best over-all fit requires $g_{\epsilon NN}^2/4\pi \approx 13.5$ with a cutoff at 0.4 F. Our conclusion is then, that the addition of the ϵ meson with a coupling constant $g_{\epsilon NN}^2/4\pi \simeq 13.5$ leads in general to improvements over our results of Paper I. A fitting of the experimental results using suitable minimization techniques could possibly make the agreement even better. Finally, a fit including a $T=1$ scalar meson could also be attempted, although it is likely that that would take us too far from the framework of Sudarshan's theory in which our calculation is based.

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Consequences of Unitarity in Some Models of CP Violation*

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Using the model-dependent assumption of 2π saturation of the unitarity sum, we find estimates for $|\eta_{00}/\eta_{+-}|$ and $\mathrm{Re}\epsilon/|\eta_{+-}|$ for a certain class of theories. Two models are tested, and $|\eta_{00}/\eta_{+-}|$ is found to be different from the originally estimated value, which was based on $\mathrm{Re}\epsilon$ as input.

The directly measured CP-violation parameters in the $K_L^0 \to 2\pi$ and $K_L^0 \to \pi l \nu$ decay experiments are η_{+-}, η_{00}, and $\mathrm{Re}\epsilon$.[1] In most theories of CP violation which have been proposed to explain the K_L^0-decay experiments, the only quantity rigorously predicted by the theory is the order of magnitude of the symmetry-breaking effect. Only the so-called superweaklike models[2] find strict values for $|\eta_{00}/\eta_{+-}|$, ϕ_{+-}, ϕ_{00}, and $\mathrm{Re}\epsilon$. When dealing with other theories, it thus appears as if much freedom is allowed for the actual values of these parameters, and estimates are sometimes based on the use of the experimental value of one of them as input ($\mathrm{Re}\epsilon$ in general).

It is the purpose of this note to show that the unitarity condition[3] can be used to provide estimates for $|\eta_{00}/\eta_{+-}|$ as well as $\mathrm{Re}\epsilon/|\eta_{+-}|$ in the framework of a certain class of CP-nonconservation models. This, of course, is in conflict with the use of $\mathrm{Re}\epsilon$ as input and leads to different results for models where that was done originally.

Following the standard phenomenological analysis of the K^0-$\bar{K}^0$ system,[4] let us denote the matrix element of the decay of K^0 into a two-pion standing-wave state with isospin I by $|A_I|e^{i\phi_I}$. The short- and long-lived neutral-kaon states are written in terms of the eigenstates of hypercharge as

$$|K_{S,L}^0\rangle = [2(1+|\epsilon_0|^2)]^{-1/2}\{(1+\epsilon_0)|K^0\rangle \pm (1-\epsilon_0)|\bar{K}^0\rangle\}. \quad (1)$$

Unlike Wu and Yang,[4] who choose $\phi_0=0$, we use the phase convention in which ϵ_0 is real. This is the phase convention in which the CP-violation phases ϕ_I are measurable quantities.[5]

Using the approximate $|\Delta I|=\frac{1}{2}$ rule for $K_S^0 \to 2\pi$ and the smallness of the observed CP-violation effect, one finds

$$\eta_{+-} \approx \epsilon_0 + i\phi_0 + (|A_2|/|A_0|\sqrt{2})\phi_2 e^{i\delta},$$
$$\eta_{00} \approx \epsilon_0 + i\phi_0 - (|A_2|\sqrt{2}/|A_0|)\phi_2 e^{i\delta}, \quad (2)$$

where $\delta = \frac{1}{2}\pi + \delta_2 - \delta_0$.

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