Distorted-field approximation to the static pion potential

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The interaction potentials for two static sources of the π-meson field are calculated in the distorted-field approximation. The spin- and isospin-dependent parts of the potential are shown to be close to the corresponding one-pion-exchange potentials. In addition, there is a very strong Wigner potential that is not present in the one-pion-exchange potential. All of the components of the potential depend significantly on the shape of the assumed pion-nucleon Yukawa interaction.

I. INTRODUCTION

The calculation of the interaction potential energy between static sources of the meson field has recently become an interesting field because of several developments. First, it seems likely that the form factor for the pion-nucleon Yukawa interaction is related to the quark-gluon structure of the nucleon core,[1] so that any measurable quantity that is sensitive to this form factor may eventually give information about the nucleon core. It is shown in the following that the static pion potential is indeed sensitive to the structure of the Yukawa interaction at the πNN vertex, so that accurate determination of the nucleon-nucleon potential may provide information about the πNN vertex. Second, the development of meson field theories of nuclear binding and structure[2] has so far not been extended to the two-nucleon problem, and the constraint of fitting the two-nucleon data has not yet been applied to the parameters of the meson theories. The static meson potential provides a way of computing two-nucleon properties from meson theory that is valid for the strong couplings that are currently fashionable. Third, recent work[3] has shown that the one-pion-exchange (OPE) potential describes several features of the two-nucleon system quite well; it is therefore important to know the extent to which the actual nucleon-nucleon potential due to pion exchange is represented by the OPE potential, and a quantitative answer to this question can be provided by the static pion potential.

An earlier paper[4] on static meson potentials developed the basic ideas needed to treat two interacting static sources of the meson field. That paper gave the first correct treatment of the static potential due to isoscalar vector-meson exchange. In the case of non-Abelian source current operators, the use of just two modes of the meson field was shown to be adequate for the calculation of the static meson potential, and a coherent-state method valid for all coupling strengths was used to treat both the single source and the two interacting sources. The asymptotic behavior of the potential was not handled correctly in Ref. 4; suitable methods for handling the asymptotic region were developed subsequently[5] and shown to give a new distorted-field approximation (DFA) scheme for computing the potential. These techniques were then supplemented by special methods developed to treat the short-range potential and used to compute the potentials in the case of isovector scalar meson field interaction,[6] in order to show that the static meson potentials in the isovector scalar case depend on the form factor, and, therefore, that there are cases in which the determination of the static meson potential can be used to gain information about the mesonic form factor of the individual sources.

The present paper applies similar methods to the case of the static pion potential, with results that are shown and discussed in Sec. VI. The general Hamiltonian for the interaction of pions with static sources is given in Sec. II. Section III summarizes the aspects of the static model for a single source of pion field that are needed for the two-source problem, and Sec. IV describes the DFA for the pion-field case. The discussion in Secs. II—IV follows the procedure in the same sections of Ref. 6, with deviations appropriate to the pseudoscalar nature of the field in the present case. Section V summarizes efforts to find an approximation that excels the DFA for small source separations. The results of the numerical calculations of the potentials are discussed in Sec. VI, and a summary of the paper is given in Sec. VII.

II. FORMULATION OF THE PROBLEM

Consider the case of two sources of pion field; the Hamiltonian is taken to be of the Yukawa form, as in Refs. 4—6,

$$
\begin{aligned}
H &= H_\omega + H_I + H_I^\dagger ,\\
H_\omega &= \int \omega(k) a_\lambda^\dagger(\mathbf{k}) a_\lambda(\mathbf{k}) d\mathbf{k} , \qquad (2.1)\\
H_I &= -\frac{f}{m}\sum_{p=1}^{2} \tau_\lambda^p \sigma_j^p \int \frac{k_j \tilde{u}^*(k)}{[16\pi^3\omega(k)]^{1/2}} a_\lambda(\mathbf{k}) e^{i\mathbf{k}\cdot\mathbf{R}_p} d\mathbf{k} ,
\end{aligned}
$$

where $\mathbf{R}_p$ is the position of the pth source; it is assumed that all sources have the same form factor $\tilde{u}(k)$ and coupling constant f. The coupling constant γ that appears later is related to the coupling constant f by

$$\gamma = \frac{f^2}{4\pi} ; \qquad (2.2)$$

for the pion case the renormalized coupling constant γ_R is 0.08. The isospin index $\lambda = 1,2,3$ and the space index $j = x,y,z$ are subject to the usual summation convention.

As in Ref. 5, single-source normalized meson mode functions $\phi_j(k)$ are defined by

$$\frac{fk_j\tilde{u}(k)}{m[16\pi^3\omega(k)]^{1/2}}=G\omega(k)\phi_j(\mathbf{k}) , \tag{2.3}$$

where the normalization constant and dimensionless coupling constant G is given by

$$G^2=\frac{f^2}{3m^2}\int\frac{k^2|\tilde{u}(k)|^2}{16\pi^3\omega^3(k)}d\mathbf{k} . \tag{2.4}$$

The absence of any extra energy parameter in the definition of ϕ follows from arguments given in Refs. 7 and 8. With the definition of ϕ of (2.3), H_I takes the form

$$H_I=-G\sum_p \tau_\lambda^p\sigma_j^p\int\omega(k)\phi_j^*(\mathbf{k})a_\lambda(\mathbf{k})e^{i\mathbf{k}\cdot\mathbf{R}_p}d\mathbf{k} . \tag{2.5}$$

III. SINGLE SOURCE

As in Refs. 5 and 6, it is necessary to solve for the single-source energy in order to compute the DFA to the two-source potential energy. When there is just one source, it can be taken to be at the origin of coordinates, so that the three normalized meson mode functions for the single source are just the functions $\phi_j(\mathbf{k})$ of (2.3); these are then used to decompose the pion-field annihilation operator $a_\lambda(\mathbf{k})$ into internal and external parts:

$$a_\lambda(\mathbf{k})=A_{\lambda j}\phi_j(\mathbf{k})+a_{\lambda\perp}(\mathbf{k}) , \tag{3.1}$$

where the $\perp$ subscript is used to indicate orthogonality to the internal mode functions $\phi_j(\mathbf{k})$;

$$\int\phi_j^*(\mathbf{k})a_{\lambda\perp}(\mathbf{k})d\mathbf{k}=0,\quad j=1,2,3 . \tag{3.2}$$

Corresponding to the decomposition (3.1) of the field, the single-source Hamiltonian also splits into internal, external, and coupling parts:

$$\begin{aligned}
&H=H_A+H_{\omega\perp}+H_1+H_1^\dagger ,\\
&H_A=W_\infty h_A ,\\
&h_A=A^\dagger\cdot A-G\rho\cdot(A^\dagger+A) ,\\
&H_{\omega\perp}=\int\omega(k)a_{\lambda\perp}^\dagger(\mathbf{k})a_{\lambda\perp}(\mathbf{k})d\mathbf{k} ,\\
&H_1=(A^\dagger-G\rho)\cdot J^{\text{ext}} ,\\
&J_{\lambda j}^{\text{ext}}=\int[\omega(k)\phi_j(\mathbf{k})]_\perp^* a_{\lambda\perp}(\mathbf{k})d\mathbf{k} ,
\end{aligned} \tag{3.3}$$

where $\rho_{\lambda j}$ stands for $\tau_\lambda\sigma_j$, and the dot product now involves both the λ and j indices; W_∞ (single source infinitely far removed from other sources) is given by

$$W_\infty=\int\omega(k)|\phi_j(\mathbf{k})|^2d\mathbf{k} . \tag{3.4}$$

It has been shown[9] that the ground-state energy of the internal Hamiltonian H_A is a good approximation to the ground-state energy of H. Reference 10 is a detailed study of the properties of the Hamiltonian h_A; in particular, it gives methods for the accurate computation of the ground-state energy of h_A for weak and intermediate values of the normalized coupling constant G.

From the commutation relation of A with H_A, it follows that if $|\alpha\mu\rangle$ is a multiplet eigenstate of H_A,

$$H_A|\alpha\mu\rangle=\epsilon_\alpha|\alpha\mu\rangle,\quad \mu=1,2,\ldots,d(\alpha) , \tag{3.5}$$

then

$$\langle\alpha\mu|A|\alpha\nu\rangle=G\langle\alpha\mu|\rho|\alpha\nu\rangle . \tag{3.6}$$

The renormalization constant $r(G)$ is obtained from the lowest isospin-$\frac{1}{2}$ spin-$\frac{1}{2}$ eigenstate $|g\rangle$ of H_A by

$$\langle g|\rho|g\rangle=r(G)\rho . \tag{3.7}$$

IV. TWO SOURCES

As in Refs. 4–6, the two-source normalized meson mode functions $\phi_{j\pm}(\mathbf{k})$ are defined by

$$\phi_{js}(\mathbf{k})=\frac{\phi_j(k)}{n_{js}\sqrt{2}}(e^{-i\mathbf{k}\cdot\mathbf{R}_1}+se^{-i\mathbf{k}\cdot\mathbf{R}_2}) , \tag{4.1}$$

where s takes the values $+$ and $-$, or sometimes $+1$ and -1. The mode normalization constants n_{js} are given by

$$\begin{aligned}
&n_{js}{}^2(\mathbf{R})=1+sc_j(\mathbf{R})\\
&c_j(\mathbf{R})=\int|\phi_j(\mathbf{k})|^2\cos\mathbf{k}\cdot\mathbf{R}\,d\mathbf{k} ,
\end{aligned} \tag{4.2}$$

and $\mathbf{R}$ is the source separation vector

$$\mathbf{R}=\mathbf{R}_1-\mathbf{R}_2 . \tag{4.3}$$

In terms of the orthonormal mode functions $\phi_{js}(\mathbf{k})$, the two-source interaction Hamiltonian H_I of (2.5) is

$$H_I=-\sum_{js}G_{js}(\mathbf{R})\rho_{\lambda js}\int\omega(k)\phi_{js}^*(\mathbf{k})a_\lambda(\mathbf{k})d\mathbf{k} , \tag{4.4}$$

where the operators $\rho_{\lambda js}$ are (note that they are defined as in Refs. 5 and 6, not as in Ref. 4)

$$\rho_{\lambda js}=\frac{1}{\sqrt{2}}(\rho_{\lambda j}^1+s\rho_{\lambda j}^2) , \tag{4.5}$$

and $G_{js}(\mathbf{R})$ is given by

$$G_{js}(\mathbf{R})=Gn_{js}(\mathbf{R}) . \tag{4.6}$$

The decomposition of the field operator $a_\lambda(\mathbf{k})$ in terms of the mode functions $\phi_{js}(\mathbf{k})$ gives

$$a_\lambda(\mathbf{k})=A_{\lambda js}\phi_{js}(\mathbf{k})+a_{\lambda\perp}(\mathbf{k}) , \tag{4.7}$$

where $a_{\lambda\perp}(\mathbf{k})$ is now orthogonal to all of the mode functions $\phi_{js}(\mathbf{k})$

$$\int\phi_{js}^*(\mathbf{k})a_{\lambda\perp}(\mathbf{k})d\mathbf{k}=0,\quad s=+,-,\quad j=x,y,z , \tag{4.8}$$

and the summation convention is now used for the sign subscript s. In the case of two sources, the subscript $\perp$ will be used generally to indicate orthogonality to the mode functions ϕ_{js}. Substitution of (4.7) into the Hamiltonian gives

$$H = H_A + H_{\omega\perp} + H_1 + H_1^\dagger ,$$
$$H_A = \sum_{js} W_{js}(\mathbf{R})[A_{\lambda js}^\dagger A_{\lambda js} - G_{js}(\mathbf{R})\rho_{\lambda js}(A_{\lambda js}^\dagger + A_{\lambda js})] , \tag{4.9}$$
$$H_{\omega\perp} = \int \omega(k) a_{\lambda\perp}^\dagger(\mathbf{k}) a_{\lambda\perp}(\mathbf{k}) d\mathbf{k} ,$$
$$H_1 = \sum_{js} [A_{\lambda js}^\dagger - G n_{js}(R)\rho_{\lambda js}] \int [\omega(k)\phi_{js}(\mathbf{k})]_\perp^* a_{\lambda\perp}(\mathbf{k}) d\mathbf{k} ,$$

where $W_{js}(\mathbf{R})$ is given by

$$W_{js}(\mathbf{R}) = \frac{1}{n_{js}^2(\mathbf{R})} \int \omega(k) \, |\phi_j(\mathbf{k})|^2 (1 + s \cos\mathbf{k}\cdot\mathbf{R}) d\mathbf{k}$$
$$= \frac{W_\infty + s w_j(\mathbf{R})}{1 + s c_j(\mathbf{R})} , \tag{4.10}$$
$$w_j(\mathbf{R}) = \int \omega(k) \, |\phi_j(\mathbf{k})|^2 \cos\mathbf{k}\cdot\mathbf{R} d\mathbf{k} .$$

W_∞ is the same as in the single-source case,

$$W_\infty = w_j(0) \tag{4.11}$$

and

$$[\omega(k)\phi_{js}(\mathbf{k})]_\perp = [\omega(k) - W_{js}(\mathbf{R})]\phi_{js}(\mathbf{k}) . \tag{4.12}$$

The term H_A in H is the "internal" part of the Hamiltonian, involving just the internal modes ϕ_{js}. The term $H_{\omega\perp}$ gives the noninteracting energy of the external modes created by $a_\perp^\dagger$, and H_1 describes the interaction between the internal and external modes. For many purposes, the ground-state energy of H_A gives a useful approximation to the ground-state energy of the total Hamiltonian; in this paper, some approximations to the ground-state energy of H_A are computed. It has been suggested[11] that the external modes be treated by diagonalizing the total Hamiltonian successively in spaces with 0,1,2,... external mesons.

The first step in dealing with H_A is, as was noted in Ref. 5, to separate the functions $W_{js}(\mathbf{R})$ and $G_{js}(\mathbf{R})W_{js}(\mathbf{R})$ into spherically symmetric s-wave and tensor d-wave parts. Let $\mathbf{R}$ be in the z direction, so that $W_{xs}(\mathbf{R}) = W_{ys}(\mathbf{R})$ and $G_{xs}(\mathbf{R}) = G_{ys}(\mathbf{R})$; then the s-wave (0) and d-wave (2) parts of W and GW are

$$W_{0s}(R) = \tfrac{1}{3}[W_{zs}(R) + 2W_{xs}(R)] ,$$
$$G_{0s}(R)W_{0s}(R) = \tfrac{1}{3}[G_{zs}(R)W_{zs}(R) + 2G_{xs}(R)W_{xs}(R)] , \tag{4.13}$$
$$W_{2s}(R) = \tfrac{1}{3}[W_{xs}(R) - W_{zs}(R)] ,$$
$$G_{2s}(R)W_{2s}(R) = \tfrac{1}{3}[G_{xs}(R)W_{xs}(R) - G_{zs}(R)W_{zs}(R)] .$$

With these definitions, H_A splits into a spherically symmetric part $H_{A0}(R)$ and a spin-2 part $H_{A2}(R)$:

$$H_A = H_{A0}(R) + H_{A2}(R) ,$$
$$H_{AL}(R) = W_{Ls}(R)\{A_s^\dagger, A_s\}^L - G_{Ls}(R)\{\rho_s, (A_s^\dagger + A_s)\}^L , \tag{4.14}$$

where, for present purposes, the $L=0$ and $L=2$ combinations are

$$\{A,B\}^0 = A\cdot B ,$$
$$\{A,B\}^2 = A_{\lambda x}B_{\lambda x} + A_{\lambda y}B_{\lambda y} - 2A_{\lambda z}B_{\lambda z} . \tag{4.15}$$

As was shown in Ref. 5, in the asymptotic region it is useful to introduce the distorted-mode functions ϕ_{jp}, $p=1,2$, that go over into the single-source mode functions as $R\to\infty$ and the corresponding single-source mode annihilation operators A_{jp}; these modes and operators are given by

$$\phi_{js}(\mathbf{k}) = \frac{1}{\sqrt{2}}[\phi_{j1}(\mathbf{k}) + s\phi_{j2}(\mathbf{k})] ,$$
$$A_{\lambda js} = \frac{1}{\sqrt{2}}(A_{\lambda j1} + sA_{\lambda j2}) . \tag{4.16}$$

In terms of these modes, H_{A0} and H_{A2} become

$$H_{AL} = \sum_p H_{AL}^p(R) + H_{ALI}(R) ,$$
$$H_{AL}^p(R) = W_L(R)\{A_p^\dagger, A_p\}^L - G_L(R)\{\rho_p, (A_p^\dagger + A_p)\}^L ,$$
$$H_{ALI}(R) = \tfrac{1}{2}[W_{L+}(R) - W_{L-}(R)](\{A_1^\dagger, A_2\}^L + \{A_2^\dagger, A_1\}^L)$$
$$- \tfrac{1}{2}[G_{L+}(R)W_{L+}(R) - G_{L-}(R)W_{L-}(R)]$$
$$\times[\{\rho_1, (A_2^\dagger + A_2)\}^L + \{\rho_2, (A_1^\dagger + A_1)\}^L] , \tag{4.17}$$
$$W_L(R) = \tfrac{1}{2}[W_{L+}(R) + W_{L-}(R)] ,$$
$$G_L(R) = \frac{G_{L+}(R)W_{L+}(R) + G_{L-}(R)W_{L-}(R)}{W_{L+}(R) + W_{L-}(R)} ,$$

where ρ_p is used interchangeably with ρ^p. The Hamiltonian $H_{A0}^p(R)$ has the same form as the single-source Hamiltonian H_A of (3.3), so that its properties are the same as those of the single-source Hamiltonian with appropriately chosen values of the parameters W_∞ and G.

The distorted-field approximation[5] (DFA) uses the product ground state

$$|g_{12}\rangle = |g_1\rangle |g_2\rangle \tag{4.18}$$

of $H_{A0}^1 + H_{A0}^2$ as an approximate ground state of the two-source system. The single-source results give

$$\langle g_p | A_p | g_p \rangle = G_0(R)\langle g_p | \rho_p | g \rangle ,$$
$$\langle g_p | \rho^p | g_p \rangle = r[G_0(R)]\rho^p = r_0(R)\rho_p , \tag{4.19}$$

so that the expectation value of $H_A(R)$ in this approximate ground state is

$$E_{df}(R) = 2E_1(R) + \sum_{L=0,2} U_L(R)\{\rho_1, \rho_2\}^L , \tag{4.20}$$

where $E_1(R)$ is the ground-state energy of $H_A{}^1(R)$ and $U_L(R)$ is given by

$$U_L(R) = r_0{}^2(R)G_0(R)\{G_0(R)[W_{L+}(R) - W_{L-}(R)] - 2[G_{L+}(R)W_{L+}(R) - G_{L-}(R)W_{L-}(R)]\} . \tag{4.21}$$

Thus, the DFA to the potential energy of the two sources is

$$V_{df}(R) = 2E_1(R) - 2E_1(\infty) + \sum_{L=0,2} U_L(R)\{\rho_1,\rho_2\}^L . \quad (4.22)$$

Since **R** is in the z direction, it follows that

$$\{\rho_1,\rho_2\}^2 = -\tau_1\cdot\tau_2(3\sigma_1\cdot\hat{\mathbf{R}}\sigma_2\cdot\hat{\mathbf{R}} - \sigma_1\cdot\sigma_2) , \quad (4.23)$$

and this gives the relation between $\{\rho_1,\rho_2\}^2$ and the usual forms of the tensor operator.

The term $2E_1(R) - 2E_1(\infty)$ is a spin- and isospin-independent or Wigner term in the potential; the other two terms are of the same type as appears in the usual one-pion-exchange calculations.

V. TREATMENT OF THE SHORT-RANGE HAMILTONIAN

The form of the Hamiltonian that is most useful for small values of the source separation is given in (4.9). In the case of the isovector scalar meson field, special methods were applied to the analogous short-range form[6] in order to get values for the potential energy that are better than the DFA values for small values of R. (Since all the calculations are variational, it is simple to decide which values are better.) Similar methods were tried for the pion-field case, but these methods did not give better values of the potential energy for interesting values of R. Therefore, only the DFA results are shown in the following section.

VI. CALCULATIONS

In order to examine the effects of changes in the form factor, the simple form

$$\tilde{u}(k) = \frac{1}{1+k^2/\Lambda^2} \quad (6.1)$$

was used for the form factor, where the parameter Λ determines the size of the source. For each value of the source-size parameter Λ, the bare coupling constant f was adjusted to make the renormalized coupling constant $f_R{}^2/4\pi$, computed by using the methods described in Ref. 10, equal to 0.08, and then the static pion potentials were computed in the DFA. The results for the two values $3m$ and $7m$ of the source-size parameter Λ are shown in Figs. 1–3 along with the corresponding one-pion-exchange (OPE) potentials computed with the renormalized coupling constant and the source density given by the form factor of Eq. (6.1).

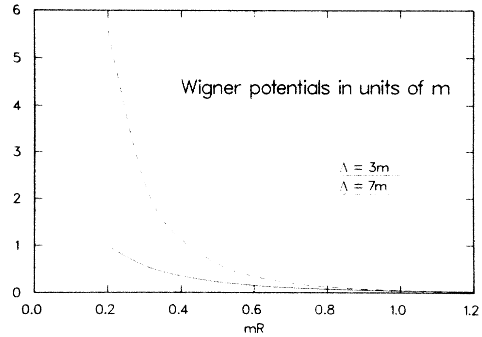

FIG. 1. The Wigner term in the static pion potential for two values of the source-size parameter Λ.

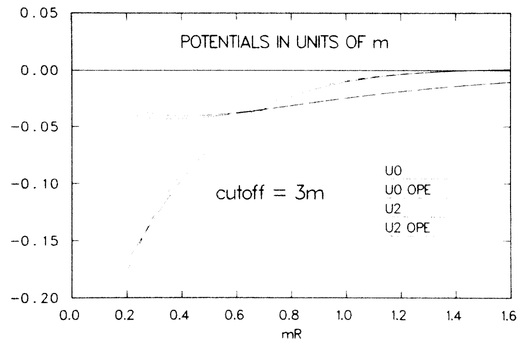

FIG. 2. The spin- and isospin-dependent terms U_0 and U_2, together with the corresponding OPE values for $\Lambda = 3m$, where Λ is the source-size parameter.

The most striking feature of these results is the rather large Wigner potential, shown in Fig. 1, that occurs in both cases. This potential is qualitatively similar to the potential energy of two Skyrmions,[12] although the Hamiltonians for the cases are entirely different. From Eq. (4.22) it is evident that this Wigner term arises from the fact that the presence of a neighboring source interferes with a particular source's optimization of its own pion field. The dependence of the Wigner term on the source-size parameter Λ is encouraging, in that it shows that an accurate determination of the Wigner term in the nucleon-nucleon potential can give information about the nature of the form factor in the pion-nucleon Yukawa interaction and, hence, about the size of the quark-gluon core of the nucleon.

In view of the strength of the pion-nucleon coupling, it is somewhat surprising that the spin- and isospin-dependent parts U_0 and U_2 of the potential, shown in Figs. 2 and 3, are very close to the corresponding OPE potentials. For these components of the potential, this means that the OPE values are probably valid to rather smaller separations than have previously been supposed; this behavior is consistent with the results of Ref. 3. These components, like the Wigner term, show a marked dependence on the shape of the pion-nucleon Yukawa interaction.

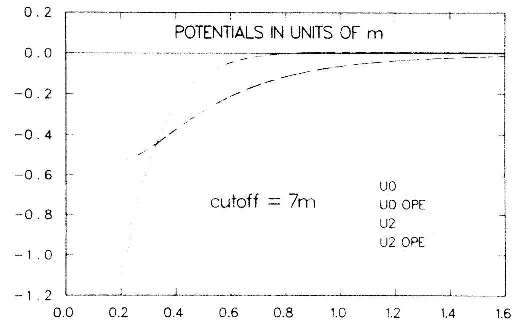

FIG. 3. The spin- and isospin-dependent terms U_0 and U_2, together with the corresponding OPE values for $\Lambda = 7m$, where Λ is the source-size parameter.

VII. SUMMARY

The potential energy of interaction of two static sources of the π-meson field has been computed in the distorted-field approximation. The spin- and isospin-dependent parts of the potential were shown to be close to the corresponding one-pion-exchange potentials. In addition, there is a strong Wigner potential that is not present in the one-pion-exchange potential. All of the components of the potential depend significantly on the shape of the assumed pion-nucleon Yukawa interaction.

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