

Short-range repulsion and isospin dependence in the kaon-nucleon (KN) system

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The short-range properties of the kaon-nucleon (KN) interaction are studied within the meson-exchange model of the Jülich group. Specifically, dynamical explanations for the phenomenological short-range repulsion, required in this model for achieving agreement with the empirical KN data, are explored. Evidence is found that contributions from the exchange of a heavy scalar-isovector meson [$a_0(980)$] as well as from genuine quark-gluon exchange processes are needed. Taking both mechanisms into account, a satisfactory description of the KN phase shifts can be obtained without resorting to phenomenological pieces.

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I. INTRODUCTION

The kaon-nucleon (KN) system provides an ideal setting for studying short-distance effects of the hadron-hadron force. This is because pions play a much less important role in the KN system than in the most extensively studied nucleon-nucleon (NN) system. Indeed, the one-pion exchange is absent in the KN interaction and the contributions from 2π exchange to the interaction seem to be weaker than in the NN system. Under such circumstances one expects that short-distance effects can be most easily isolated from the attractive medium-range background and that possibly effects from explicit quark-gluon degrees of freedom can be identified.

A large body of work accumulated in the last 50 years indicates that meson degrees of freedom are very efficient for describing low-energy hadron-hadron interactions. In particular, for the KN system, a few years ago the Jülich group presented a meson-exchange model for the K^+N scattering [1,2]. Reference [1] considered single boson exchanges (σ, ρ, ω), together with contributions from higher-order diagrams involving N , Δ , K , and K^* intermediate states. It turned out that the S -wave observables of KN experiments could only be described with the model if the value of the $KK\omega$ coupling constant is increased about 60% above the value that follows from the SU(3) (quark flavor) symmetry. Specifically, the increased value of the $KK\omega$ coupling constant was strictly necessary to obtain the S -wave low-energy parameters and the energy dependence of S -wave phase shifts for both isospin $I=0$ and $I=1$ channels. However, this increased value leads to additional repulsion in the P and higher partial waves which seemed to be not favored by the empirical data, especially in the P_{03} and P_{13} channels.¹ Thus, it was concluded in Ref. [1] that the required additional contributions must be much shorter ranged than the ω exchange.

Further evidence for the conjecture that the repulsion needed to describe KN scattering cannot be interpreted com-

pletely in terms of conventional ω exchange came from subsequent investigations of the $\bar{K}N$ system [3]. In a meson-exchange model like the one developed by the Jülich group there is a close connection between the KN and $\bar{K}N$ interactions due to G -parity conservation. Specifically, this means that the repulsive ω exchange changes sign for K^-N , because of the negative G parity of the ω meson, and becomes attractive. A large contribution from the ω exchange as favored by the KN S waves turns then into a strongly attractive piece—which is indeed much too strong to fit the K^-N data [3].

The conclusion from those results was that ω exchange, as treated in this model, can only be interpreted as an effective contribution that parametrizes besides the “physical” ω exchange also further shorter-ranged mesonic contributions or genuine quark-gluon effects or both. This was shown by a model analysis where the coupling constants of the ω meson ($g_{KK\omega}$, $g_{NN\omega}$) were kept at their SU(3) symmetry values and an additional phenomenological (extremely short-ranged) repulsive contribution, a “ σ_{rep} ,” with a mass of about 1.2 GeV was added—see Figs. 1(a) and 1(b).

In Ref. [2] the model was further refined by replacing σ and ρ exchange by the correlated 2π -exchange contribution in the $J^P=0^+$ and $J^P=1^-$ channels, respectively, as illustrated in Fig. 1(c). This was done by starting from a microscopic model for the t -channel reaction $N\bar{N}\rightarrow K\bar{K}$ with $\pi\pi$ (and $K\bar{K}$) intermediate states and using a dispersion relation over the unitarity cut. Such a realistic model of (effective) σ and ρ exchange was then used to reconstruct an extended meson-exchange model for KN scattering. But again, as in Ref. [1], the addition of a phenomenological σ_{rep} was essential to describe the data with the SU(3) $KK\omega$ coupling constant.

One possible interpretation for the need of a very short ranged repulsion, shorter ranged than that provided by ω exchange, is that quark-gluon effects are playing a role [1]. The study of the KN interaction in the context of quark models has a long history since the 1980s [4]. More recently, the subject has gained renewed interest with the works of Barnes and Swanson [5] and Silvestre-Brac and collaborators [6–8]. The main ingredients in the calculations of both groups are

¹The spectroscopic notation used is such that a partial wave with angular momentum L , total angular momentum J , and isospin I is denoted by L_{12J} .

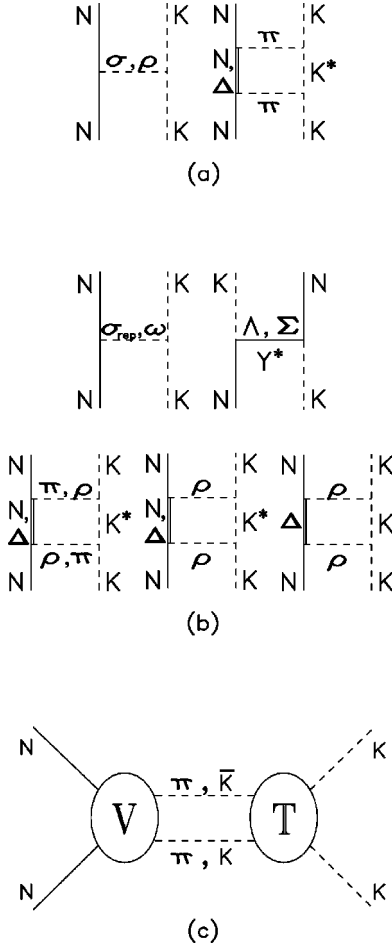


FIG. 1. Meson-exchange contributions to KN scattering in the Jülich model [1,2]. Diagrams (a) and (b) define the model of Ref. [1] and diagram (c) is the correlated 2π exchange calculated in Ref. [2], which was parametrized by diagram (a) in Ref. [1].

the nonrelativistic quark model and a quark interchange mechanism with one-gluon exchange (OGE). One important conclusion of these calculations is that the derived KN interaction is short ranged and repulsive, and strongly isospin dependent. As we discussed above, although in the Jülich model the overall strength and energy dependence of the S wave phase shift of K^+N scattering can be obtained by augmenting the value of the $KK\omega$ coupling, the P and higher partial waves do not come out right and the introduction of the exchange of a fictitious scalar particle with repulsive character was essential for this matter. In view of this, the substitution of an isospin independent σ_{rep} by a strongly isospin-dependent quark-gluon dynamics is not trivial and apparently bound to fail. However, we note that in both Refs. [1,2] the $a_0(980)$ meson was left out without any apparent reason. This meson, being a scalar isovector, is an important source of isospin dependence and it has a mass not much larger than those of the other mesons considered in the model. Thus, it can, in principle, play an important role in the isospin dependence of the KN interaction.

The main motivation of our paper is to investigate the energy and isospin dependences of the KN interaction in a hybrid model, in which the Jülich model is extended by add-

ing the $a_0(980)$ exchange, and where the very short ranged part of the KN interaction is described by quark-gluon exchange, instead of the phenomenological σ_{rep} . We construct an effective KN potential from a microscopic nonrelativistic quark model from the interquark exchange mechanism and use all components of the OGE interaction. These include the Coulomb, spin-independent contact, and spin-orbit interactions. In addition, we use a linearly rising potential to represent quark confinement.

The paper is organized as follows. In the following section we describe the KN meson-exchange model of the Jülich group. In Sec. III we provide an overview of studies on the KN system that were carried out in the framework of the quark model. In Sec. IV we present and discuss our results. Specifically, we investigate the consequences of replacing the phenomenological σ_{rep} of the Jülich model by quark-gluon exchange. Our conclusions and perspectives are presented in Sec. V. The Appendix presents the expressions for the effective KN potentials in the quark model.

II. THE JÜLICH KN MODEL

The Jülich meson-exchange model of the KN interaction has been widely described in the literature [1,2,9,10] and we refer the reader to those works for details. Here we will only summarize the features that are relevant to the present study.

The Jülich meson-exchange model of the KN interaction was constructed along the lines of the (full) Bonn NN model [11] and its extension to the hyperon-nucleon (YN) system [12]. Specifically, this means that one has used the same scheme (time-ordered perturbation theory), the same type of processes, and vertex parameters (coupling constants, cutoff masses of the vertex form factors) fixed already by the study of these other reactions.

The diagrams considered for the KN interaction are shown in Fig. 1. Based on these diagrams a KN potential V is derived, and the corresponding reaction amplitude T is then obtained by solving a Lippmann-Schwinger type equation defined by the time-ordered perturbation theory,

$$T = V + VG_0T. \quad (1)$$

From this amplitude, phase shifts and observables (cross sections, polarizations) can be obtained in the usual way.

As seen in Fig. 1, obviously the Jülich model contains not only single-meson (and baryon) exchanges, but also higher-order box diagrams involving NK^* , ΔK , and ΔK^* intermediate states. Most vertex parameters involving the nucleon and the $\Delta(1232)$ isobar can be taken over from the (full) Bonn NN potential. The coupling constants at vertices involving strange baryons are fixed from the YN model (model B of Ref. [12]). Those quantities (g_{NAK} , $g_{N\Delta K}$, g_{NY^*K}) have been related to the empirical $NN\pi$ coupling by the assumption of SU(6) symmetry, cf. Refs. [1,2].

For the vertices involving mesons only, most coupling constants have been fixed by SU(3) relating them to the empirical $\rho \rightarrow 2\pi$ decay. Exceptions are the coupling constants $g_{KK\sigma}$ and $g_{KK\omega}$, which have been adjusted to the KN data for the following reason: The σ meson (with a mass of about

600 MeV) is not considered as a genuine particle but as a simple parametrization of correlated 2π -exchange processes in the scalar-isoscalar channel. Therefore, its coupling strength cannot be taken from symmetry relations. Concerning the ω exchange it was found that a much larger strength than obtained from SU(3) was required in order to obtain sufficient short-range repulsion for a reasonable description of the S -wave KN phase shifts [1]. The ω coupling $g_{KK\omega}$ had to be increased by about 60% over the symmetry value—quite analogous to the situation in the NN system [11]. However, such an increased ω exchange turned out to be in contradiction with the empirical data on P - and higher partial waves. Specifically, P_{03} and P_{13} do not really demand additional repulsive contributions. Thus, it was concluded that the additional repulsion should be of rather short-ranged nature. Such a contribution would still allow to obtain a reasonable description of the S waves, but would leave P and higher partial waves basically unchanged.

As a consequence, the Jülich group presented a model where the coupling strengths for both $g_{NN\omega}$ and $g_{KK\omega}$ were kept at their SU(6) values. At the same time, a phenomenological, very short ranged contribution was added. This phenomenological piece has the same analytical form as σ exchange, but an exchange mass of 1200 MeV and, most importantly, an opposite sign. Accordingly, it was denoted σ_{rep} .

In a subsequent investigation the $\sigma(600)$ and also the elementary ρ were replaced by a microscopic model for correlated 2π (and $K\bar{K}$) exchange between a kaon and a nucleon, in the corresponding scalar-isoscalar and vector-isovector channels [2]. The starting point for this was a model for the reaction $N\bar{N} \rightarrow K\bar{K}$ with intermediate 2π and $K\bar{K}$ states, based on a transition in terms of baryon ($N, \Delta, \Lambda, \Sigma$) exchange and a realistic coupled channel $\pi\pi \rightarrow \pi\pi$, $\pi\pi \rightarrow K\bar{K}$, and $K\bar{K} \rightarrow K\bar{K}$ amplitude. The contribution in the s channel is then obtained by performing a dispersion relation over the unitarity cut. But also in this model the phenomenological short-ranged σ_{rep} was needed in order to achieve agreement with the empirical phase shifts.

Since the results of Ref. [2] indicate that the contributions of the correlated 2π exchange in the scalar-isoscalar channel are in rough agreement with the effective description by σ -exchange used in Ref. [1], we will employ the latter in the present investigation for simplicity reasons. Specifically, we will use the KN model I presented in Ref. [2]. The parameters of this model are summarized in Table I. Resulting phase shifts for the Jülich model I [2] will be shown and compared with empirical data in Sec. IV. Further results, including scattering observables, can be found in Ref. [2].

III. THE KN INTERACTION IN THE QUARK MODEL

In this section we briefly review the salient features of model calculations of the KN interaction that are based on quark-gluon exchange and derived within the nonrelativistic quark model. Specifically, we will focus on the more recent calculations of Barnes and Swanson [5] and Silvestre-Brac and collaborators [6–8].

TABLE I. Vertex parameters used in the Jülich KN model I [2]. Numbers in parentheses denote corresponding values of the model discussed in the present paper, when different.

Process	Exchanged particle	M_r or m_r^a (MeV)	$g_1 g_2 / 4\pi^b$ [f_1 / g_1]	Λ_1^c (GeV)	Λ_2^c (GeV)	
$KN \rightarrow KN$	σ	600	1.300 (1.000)	1.7	1.5 (1.2)	
	σ_{rep}	1200	-40 (\cdot)	1.5 (\cdot)	1.5 (\cdot)	
	a_0	980	- (2.600)	- (1.5)	- (1.5)	
	ω	782.6	2.318 [0]	1.5	1.5	
	ρ	769	0.773 [6.1]	1.4	1.6	
	Λ	1116	0.905	4.1	4.1	
$KN \rightarrow K^*N$	Σ	1193	0.031	4.1	4.1	
	Y^*	1385	0.037	1.8	1.8	
	π	138.03	3.197	1.3	0.8	
	ρ	769	0.773[6.1]	1.4	1.0	
	$KN \rightarrow K^*\Delta$	π	138.03	0.506	1.2	0.8
		ρ	769	4.839	1.3	1.0
$KN \rightarrow K\Delta$	ρ	769	4.839	1.3	1.6	

^aMass of exchanged particle.

^bProduct of coupling constants (ratio of tensor to vector coupling).

^cCutoff mass.

Barnes and Swanson use the quark-Born-diagram (QBD) method [13]. In this method, the KN scattering amplitude is assumed to be the coherent sum of all OGE interactions followed by all allowed quark line exchanges. The input to this method is the microscopic quark-quark interaction and kaon and nucleon wave functions. Barnes and Swanson [5] used the contact spin-spin “color-hyperfine” component of the OGE and used Gaussian wave functions for the interacting hadrons, which allowed them to evaluate the KN scattering amplitude analytically. They calculated isospin $I=0$ and $I=1$ scattering observables such as S -wave phase shifts and scattering lengths. The model has only two parameters, the ratio of the u, d to s quark masses, $\rho = m_q / m_s$, and α_s / m_q^2 , where α_s is the quark-gluon coupling constant. By using typical quark-model parameters and using the Born approximation, Barnes and Swanson obtained very reasonable results for the S -wave phase shifts [5].

The studies of Silvestre-Brack and collaborators [6–8] complement the work of Barnes and Swanson in several aspects. First of all, Refs. [6–8] employ the resonating group method (RGM) instead of the QBD method. The main difference between the two methods refers to the way orthogonality effects in the relative KN wave functions are treated. These are effects due to the Pauli principle for quarks in different clusters. It can be shown [14] that the differences between effective hadron-hadron interactions calculated with both methods are usually small, although not entirely negligible. In Ref. [6] the authors calculate S -wave scattering phase shifts using the Coulomb, spin-spin, and constant contact pieces of the OGE quark-quark interaction. In addition, they include the exchange of π and σ mesons—considered

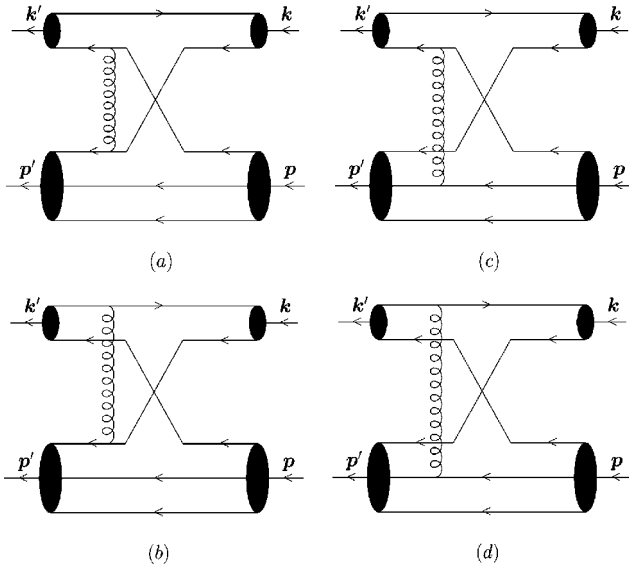


FIG. 2. The four quark-interchange kaon-nucleon scattering diagrams.

as elementary particles—between quarks, and a linearly rising confining potential. The same quark-quark interaction is used to build the K and N wave functions and to generate the KN interaction. The parameters are constrained to reproduce the low-lying meson and baryon spectrum. The main conclusion of this study was that it is impossible to describe both $I=0$ and $I=1$ isospin channels simultaneously within the model. Relativistic kinetic energy effects were investigated in Ref. [7]. The results obtained for the S waves, for which one expects such effects to be more important, were not much different from the corresponding nonrelativistic ones. In Ref. [8], scattering phase shifts from S up to G waves were calculated. The difference from Ref. [6] is that a spin-orbit interaction was added to the OGE and confining pieces used there. No meson exchanges between quarks were considered. The parameters were again fixed by requiring a good description of the low-lying meson and baryon spectrum. The results obtained are such that the $I=0$ S -wave phase shift is reasonably well described, while the corresponding $I=1$ is too repulsive. The P and higher partial-wave phase shifts are poorly described, with the exception of the P_{11} , D_{13} , D_{15} , and G_{19} phases. The P_{01} phase, for example, is predicted to be almost zero, while the corresponding experimental phase grows from zero up to 60° at $p_{lab}=1$ GeV.

The results of Silvestre-Brac and collaborators [6–8] clearly indicate that the quark-interchange mechanism with OGE alone is not sufficient to describe the K^+N data. However, it seems to provide at least enough strength for S waves. In view of the discussion above on the Jülich model, we investigate here the substitution of σ_{rep} in that model by the quark-interchange mechanism with OGE. Recently, one of us [15] has derived the contribution of the spin-spin part of the OGE to the central part of the KN interaction using the mapping formalism developed in Ref. [14]. The different contributions to the K^+N effective potential from the quark-interchange mechanism are illustrated in Fig. 2. The on-shell KN amplitude is identical to the one derived by Barnes and

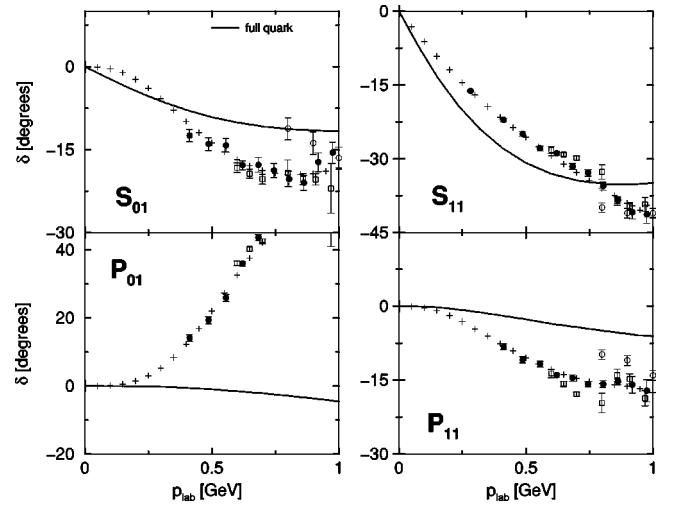


FIG. 3. KN phase shifts. The solid lines are the result from the full quark-model calculation described in the text. Experimental phase shifts are taken from Refs. [16] (open circles), [17] (open squares), and [18] (filled circles and pluses).

Swanson [5]. In order to iterate the potential in a Lippmann-Schwinger equation one needs the off-shell amplitude. For this purpose, we have calculated the contributions of all remaining components of the OGE to the K^+N effective potential within the framework of Ref. [14]. We found that the spin-spin component of the OGE gives by far the most important contribution to the K^+N effective potential.

For illustrative purposes we present, in this section, the phases calculated in Born approximation—in the following section the OGE KN potential is iterated in the Lippmann-Schwinger type of equation, Eq. (1). In Fig. 3 we present the S - and P -wave phase shifts resulting from the OGE and the confining interaction. The higher partial waves are very small and are not shown. The experimental data points in this figure are taken from Refs. [16–18]. The analytical expressions for the KN potential are given in the Appendix. The parameters of the potential are the masses of the constituent quarks, $m_q (=m_u=m_d)$ and m_s , the quark-gluon hyperfine coupling α_s , and the size parameters of the nucleon and kaon wave functions, α and β . We use the “reference parameter set” of Barnes and Swanson [5], which contains conventional quark model parameters. These are

$$\rho = m_q/m_s = 0.33 \text{ GeV}/0.55 \text{ GeV} = 0.6,$$

$$\alpha_s/m_q^2 = 0.6/(0.33)^2 \text{ GeV}^2,$$

$$\alpha = 0.4 \text{ GeV}, \quad \beta = 0.35 \text{ GeV}. \quad (2)$$

The string tension of the confining potential is taken to be $\sigma = 0.18 \text{ GeV}^2$ [19]. In addition, when calculating the phase shifts we use the physical masses of the nucleon and the kaon, $M_N = 0.940 \text{ GeV}$ and $M_K = 0.495 \text{ GeV}$. Figure 3 shows that S waves are reasonably well described by the model, although the $I=0$ phase agrees less well with the data at higher energies. The fit can be improved slightly by choosing another set of parameters, as done by Barnes and Swan-

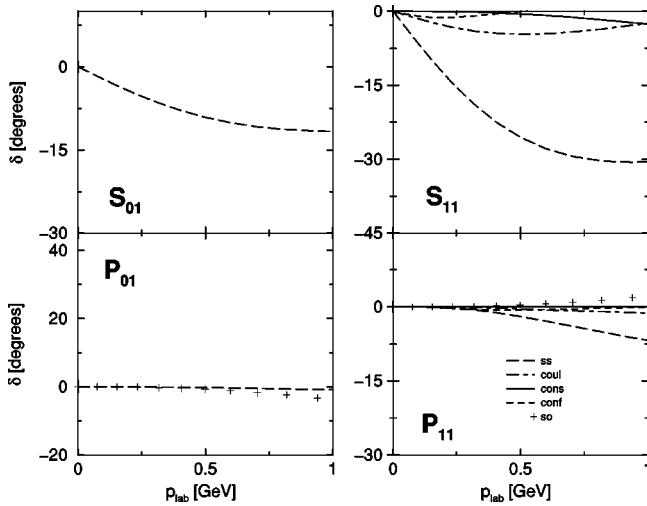


FIG. 4. KN phase shifts resulting from various contributions of the quark model, cf. Appendix: Coulomb (dash-dotted line); spin orbit (pluses); confinement (short-dashed line); contact spin-spin (long-dashed line); contact constant (solid line).

son in their study of the S -wave phase shifts [5]. In this paper we maintain the reference set, since the general trend of the higher partial waves will not be modified by a change of the quark model parameters.

In Fig. 4 we show the separate contributions of the OGE and confining interaction to S and P phases. The dominance of the spin-spin component of OGE is clearly seen. The confining interaction gives a very small contribution to all waves, and the only noticeable effect from the other components of the OGE is the one from the Coulomb part in the S_{11} wave.

IV. RESULTS AND DISCUSSION

As discussed in the last two sections, the KN interaction in the low-energy region ($p_{lab} \leq 1$ GeV/ c) can be well understood within the meson-exchange picture. However, a good quantitative overall description of the data can only be achieved by adding a phenomenological contribution that is extremely short ranged and repulsive—and therefore affects essentially the S waves only. At the same time, interaction models based on quark-gluon degrees of freedom yield only a mediocre overall reproduction of the KN phase shifts. However, the predicted S waves are in fairly good agreement with empirical results, suggesting that at least the short-ranged part of the KN interaction is well accounted for by the one-gluon-exchange mechanism that is the dominant ingredient in those quark models. It is, therefore, tempting to combine the contributions of those two complementary approaches to the KN force. Indeed, such a procedure is in the spirit of the original KN model of the Jülich group where it was suggested that the short-ranged phenomenological piece added in this model might be an effective parametrization of either further short-ranged mesonic contributions or of genuine quark-gluon effects or both [1,20].

Results for the KN phase shifts of the original Jülich model (note that we use here model I of Ref. [2]) are shown

by the dash-dotted lines in Fig. 5. If we switch off the contribution from the phenomenological σ_{rep} and add the contribution from one-gluon exchange instead we obtain the short-dashed lines. The parameters of the quark model are the same as in the preceding section. We see that the OGE is indeed capable of producing repulsive contributions which are of a comparable order of magnitude as the one of the phenomenological σ_{rep} . Indeed, after the discussion in Sec. III this could have been expected. However, it is definitely surprising that for the S_{11} partial wave the results with OGE (and without σ_{rep}) are almost identical to the ones of the original Jülich model. In case of the S_{01} , the situation is somewhat less satisfying. Here the repulsion provided by the OGE is significantly smaller than the one parametrized by the σ_{rep} . This is simply a consequence of the isospin dependence inherent in the OGE—the phenomenological σ_{rep} is, of course, per construction an isoscalar. The higher partial waves (in the $I=0$ as well as the $I=1$) are again only marginally changed as compared to the original results—testifying that also the OGE is of rather short range.

The above results can be seen as an indication that the OGE is not the only short-range physics that is parametrized by the σ_{rep} of the Jülich KN model. (Indeed, one might argue that this could have been already guessed from the difference in the isospin structure.) Besides possible higher-order contributions resulting from quark-gluon dynamics, one should not forget to take into consideration further shorter-ranged mesonic contributions. Indeed, the exchange of the $a_0(980)$ meson, which is a scalar-isovector particle, is a natural candidate for this. With its mass of about 1 GeV its contributions are definitely of short-ranged nature as required. Furthermore, its isospin structure leads to attractive contributions in the $I=1$ channel but to repulsive contributions in the $I=0$ channel. Thus, it complements the isospin dependence of the OGE in an almost ideal way and in conjunction with the latter would lead to contributions that are almost isospin independent—as those of the phenomenological σ_{rep} . The $a_0(980)$ meson is taken into account in the Bonn NN model [11] (it is denoted as δ meson there). However, for unexplained reasons, it was not included in the original Jülich KN model. Subsequent investigations of the Jülich group on the structure of the $a_0(980)$ meson suggested that this resonance can be understood in terms of strong correlations in the $\pi\eta-K\bar{K}$ channel [21]. Thus, the situation is similar to the strong $\pi\pi-K\bar{K}$ correlations in the scalar-isoscalar channel that are usually effectively parametrized by the σ meson. Accordingly, it is in the spirit of the Bonn/Jülich models of hadronic reactions to consider the contributions of the $a_0(980)$ meson, and indeed it is included in the more recently developed models of the πN [22,23] and YN [24] interactions. Following the arguments in Refs. [21,22] we do not view the a_0 exchange as a genuine meson exchange but rather as a parametrization of correlations in the mesonic systems in the scalar-isovector channel. Therefore we consider the a_0 coupling constants as free parameters. In principle, $g_{K\bar{K}a_0}$ could be determined from the a_0 decay width into the $K\bar{K}$ system. However, the experimental information on this quantity is still very poor, cf. Ref. [25],

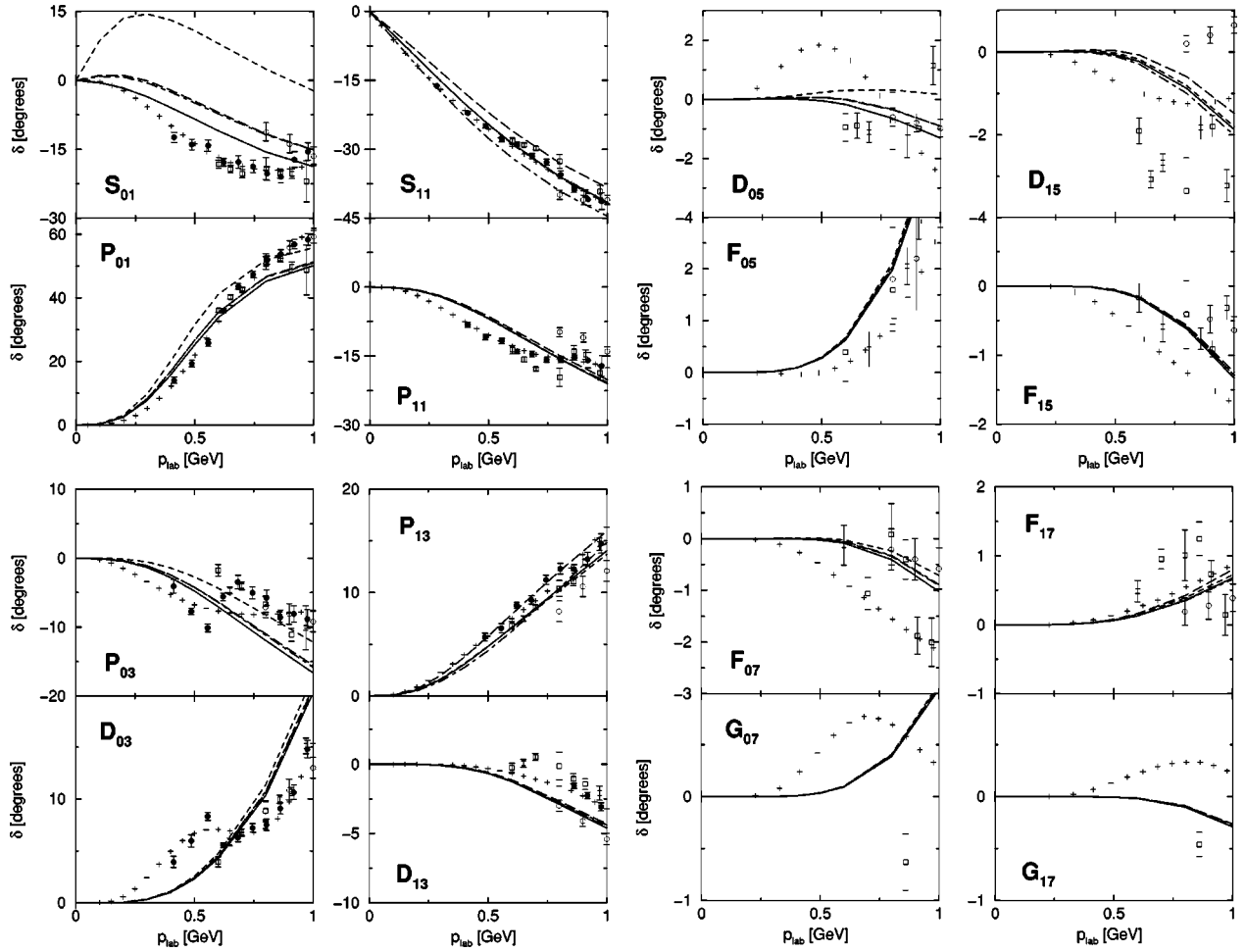


FIG. 5. KN phase shifts. The dash-dotted lines are the phase shifts of the original Jülich model I from Ref. [2]. The short-dashed line shows results where the phenomenological σ_{rep} in the Jülich model is replaced by the quark-model contribution. Adding the a_0 -exchange contribution yields the long-dashed line. The solid line is obtained after refitting the parameters of the σ . Same description of experimental phase shifts as in Fig. 3.

and thus cannot provide more than a guideline. Note that in the a_0 exchange the product of the NNa_0 and $K\bar{K}a_0$ coupling constants appears, and therefore we list only this product in Table I.

Results including now a_0 meson exchange as well as the OGE contributions are shown by the long-dashed lines in Fig. 5. The coupling strength of the a_0 exchange has been chosen in such a way that the model prediction for the S_{01} partial wave agrees roughly with the result of the original Jülich KN model. As can be seen in Fig. 5, the inclusion of the a_0 exchange influences also the KN P waves in the $I=0$ channel, i.e., the P_{01} and P_{03} , whereas all other higher partial waves remain basically unchanged. As a matter of fact, the present model based on a_0 meson exchange and OGE contributions yields pretty much the same results as the original Jülich model utilizing the phenomenological σ_{rep} . Minor differences occur only in the S_{11} partial wave, which turns out to be now somewhat too less repulsive in comparison to the data. Thus, as a last step, we have slightly readjusted the parameters of the σ meson, cf. Table I, which then leads to the final results shown by the solid lines in Fig. 5. Those results provide clear evidence that a comparable quan-

titative description of the KN interaction can be achieved with a model that avoids phenomenological contributions like the σ_{rep} of the original Jülich model.

Finally, we want to address the question whether a description of the KN interaction is possible within the Jülich model without introducing explicit contributions from the OGE. Treating the coupling constants of the σ and a_0 mesons as completely free parameters we were indeed able to obtain a reasonable reproduction of the S_{01} and S_{11} partial waves. However, it could only be achieved by assuming that the σ -exchange contribution is basically zero. Of course, this is completely unrealistic in view of the results obtained for the strength of the correlated two-pion exchange in the σ channel in Ref. [2]. Moreover, the description of the higher partial waves deteriorates significantly if the σ -exchange contribution is so strongly reduced.

V. CONCLUSIONS AND PERSPECTIVES

In this paper we have studied the short-range properties of the KN interaction. In particular, we have taken the KN meson-exchange model of the Jülich group and we explored

possible dynamical explanations for a phenomenological (extremely short-ranged) repulsive contribution, a “ σ_{rep} ” with a mass of about 1.2 GeV, that is present in the Jülich model. Such a phenomenological, repulsive, and rather short-ranged piece had to be introduced in that model for achieving agreement with the empirical KN data.

The very short-ranged nature of this repulsion could be a sign that quark-gluon dynamics is playing a role. Therefore, we have calculated corresponding contributions to the KN interaction based on the nonrelativistic quark model and a quark interchange mechanism with one-gluon exchange. It turned out that those processes are indeed short ranged and repulsive. However, unlike the phenomenological “ σ_{rep} ” in the Jülich model, they are also strongly isospin dependent. Thus, one-gluon exchange alone can certainly not explain the required short-range physics. Consequently, we examined additional short-range physics that arises in the mesonic sector, and specifically the exchange of the (scalar-isovector) $a_0(980)$ meson. Its contribution was not included in the original Jülich KN model.

Due to its isospin structure the $a_0(980)$ exchange provides attraction in the $I=1$ channel and repulsion in the $I=0$ channel, and therefore counterbalances the isospin dependence of the one-gluon exchange. Taking both mechanisms (a_0 as well as one-gluon exchange) into account yields a short ranged and repulsive but basically isospin-independent interaction—similar to the one parametrized by the σ_{rep} —and, consequently, a satisfactory description of the KN phase shifts can be obtained without resorting to phenomenological pieces, as demonstrated in the present paper.

The authors of the original Jülich KN model conjectured that the introduced phenomenological σ_{rep} might be an effective parametrization of either further short-ranged mesonic contributions or genuine quark-gluon effects or both [1]. Our investigation provides strong evidence that the third alternative is realized. Specifically, it lends support to the supposition that effects from quark-gluon degrees of freedom can be explicitly seen in the KN system. Still one can raise the question whether contributions from genuine quark-gluon dynamics are really needed. For example, could their role be taken over by the exchange of heavier vector mesons, say? To answer this question it will be very instructive to study the $\bar{K}N$ system again, using the present model. For example, the investigations in Ref. [3] have shown that the $\bar{K}N$ data require only a strongly reduced short-ranged repulsive piece, i.e., only about 20% of the phenomenological σ_{rep} used in the KN system. It will be interesting to see whether the present scenario of combined $a_0(980)$ exchange and quark-gluon dynamics is able to generate these properties when going over to the $\bar{K}N$ system. One should note, however, that the treatment of the $\bar{K}N$ channel in the nonrelativistic quark model is more complicated than the KN system since it involves s -channel gluon exchange. Special care must be taken with such processes because it is not clear that the use of perturbative massless gluons makes physical sense in this model. Contributions of intermediate hybrid $q\bar{q}g$ states to the process must certainly be considered. Still the extension of the quark interchange model to incorporate

gluon annihilation in the $\bar{K}N$ system would be a very interesting new development. Investigations along this line are planned for the future.

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APPENDIX: CONTRIBUTIONS TO THE KN INTERACTION FROM OGE AND CONFINEMENT

In this appendix we present the different contributions the OGE and from the confining potential to the KN potential used in the present paper. Let us denote the initial and final three momenta of the interacting quarks by $\vec{k}_1, \vec{k}_2, \vec{k}'_1$, and \vec{k}'_2 (prime means final state). It is convenient to define the following combinations of momenta: $\vec{q} = \vec{k}'_1 - \vec{k}_1 = \vec{k}_2 - \vec{k}'_2$, $p_1 = (\vec{k}_1 + \vec{k}'_1)/2$ and $p_2 = (\vec{k}_2 + \vec{k}'_2)/2$. In terms of these, the interquark interaction can be written as

$$H_{OGE} = \sum_{ij} \left[\sum_a \mathcal{F}^a(i) \mathcal{F}^a(j) \right] V_{ij}(\vec{q}, \vec{p}_i, \vec{p}_j), \quad (\text{A1})$$

where i, j identify the quarks (or antiquarks) 1,2 and $V_{ij}(\vec{q}, \vec{p}_i, \vec{p}_j)$ depends on spin variables and the indicated momenta. The color SU(3) matrices $\mathcal{F}^a(i)$, $a=1, \dots, 8$, are given in terms of the Gell-Mann matrices λ^a as $\mathcal{F}^a(i) = \lambda^a/2$ when i is a quark and $\mathcal{F}^a(i) = -\lambda^{aT}/2$ when i is an antiquark (T means transpose). We refer the reader to the literature for the explicit expression of $V_{ij}(\vec{q}, \vec{p}_i, \vec{p}_j)$ —see, for example, Eq. (3) of Ref. [19].

Next, we present the individual contributions in $V_{ij}(\vec{q}, \vec{p}_i, \vec{p}_j)$ to the KN potential. We represent each contribution to the KN potential as

$$V(\vec{p}, \vec{p}') = \frac{1}{2} \sum_{D=a}^d [V_D(\vec{p}, \vec{p}') + V_D(\vec{p}', \vec{p})], \quad (\text{A2})$$

where \vec{p} and \vec{p}' are the initial and final center of mass momenta of the the KN system, and the index D identifies the diagrams a, \dots, d of Fig. 2. The explicit evaluation of these diagrams requires also the specification of the nucleon and kaon wave functions, Ψ_N and Ψ_K . These are taken to be in momentum space of the form

$$\Psi_N(\vec{p}) = \delta(\vec{p} - \vec{k}_1 - \vec{k}_2 - \vec{k}_3) N(\vec{p}) \phi(\vec{k}_1) \phi(\vec{k}_2) \phi(\vec{k}_3), \quad (\text{A3})$$

where

$$\phi(\vec{k}) = \left(\frac{1}{\pi\alpha^2} \right)^{3/4} \exp(-\vec{k}^2/2\alpha^2),$$

$$N(\vec{p}) = (3\pi\alpha^2)^{3/4} \exp t(p^2/6\alpha^2), \quad (\text{A4})$$

and

$$\Psi_K(\vec{p}) = \delta(\vec{p} - \vec{k}_q - \vec{k}_{\bar{q}}) \left(\frac{1}{\pi\beta^2} \right)^{3/4} \exp \left[- \frac{(m_1 \vec{k}_q - m_2 \vec{k}_{\bar{q}})^2}{8\beta^2} \right], \quad (\text{A5})$$

with

$$m_1 = \frac{2m_{\bar{q}}}{m_q + m_{\bar{q}}}, \quad m_2 = \frac{2m_q}{m_q + m_{\bar{q}}}. \quad (\text{A6})$$

For convenience, we also introduce the quantities

$$\rho = m_q/m_s, \quad g^2 = \alpha/\beta, \quad b = 1/\alpha. \quad (\text{A7})$$

The explicit contributions of the different pieces of the OGE and of the confining potential to the effective KN interaction are given by the following expressions.

1. Coulomb

$$\begin{aligned} V_D^{Coul}(\vec{p}, \vec{p}') &= 4\pi\alpha_s \omega_D(I) \int_0^\infty ds \eta_D(s) \\ &\times \exp[-A_D(s)p^2 - B_D(s)p'^2 \\ &+ C_D(s)\vec{p} \cdot \vec{p}'], \end{aligned} \quad (\text{A8})$$

where the variable s comes in because we have chosen to perform a Laplace transform of the Coulomb potential $1/q^2$ in order to integrate over the variable q , and the coefficients $\omega_D(I)$ (I identifies isospin $I=1$ or $I=0$) come from summing over color-spin flavor of quarks. The functions A_D, B_D, C_D , and η_D for each diagram can be written as a ratio $A_D = n(A_D)/d(A_D), \dots, \eta_D = n(\eta_D)/d(\eta_D)$.

Diagram (a).

$$\begin{aligned} n(A_a) &= 24b^4\beta^2 + 9b^2m_1^2 + (80b^2\beta^2 + 24b^4\beta^4 \\ &- 48b^2\beta^2m_1 + 6m_1^2 + 18b^2\beta^2m_1^2)s, \end{aligned}$$

$$n(B_a) = n(A_a),$$

$$\begin{aligned} n(C_a) &= 8b^4\beta^2 + 3b^2m_1^2 + (-16b^2\beta^2 + 8b^4\beta^4 \\ &+ 16b^2\beta^2m_1 + 2m_1^2)s, \end{aligned}$$

$$n(\eta_a) = 3\sqrt{3}b^3/8\pi^3,$$

$$d(A_a) = d(B_a) = 6d(C_a) = 48\beta^2[3b^2 + (2 + 3b^2\beta^2)s],$$

$$d(\eta_a) = [3b^2 + (2 + 3b^2\beta^2)s]^{3/2},$$

$$\omega_a(1) = -4/9, \quad \omega_a(0) = 0. \quad (\text{A9})$$

Diagram (b).

$$\begin{aligned} n(A_b) &= -24 - 80b^2\beta^2 + 24b^4\beta^4 + 12m_1 + 60b^2\beta^2m_1 \\ &- 9b^2\beta^2m_1^2 - (96\beta^2 + 176b^2\beta^4 - 48b^4\beta^6 \\ &- 48\beta^2m_1 - 6\beta^2m_1^2 + 120b^2\beta^4m_1 - 18b^2\beta^4m_1^2)s, \end{aligned}$$

$$\begin{aligned} n(B_b) &= 12 + 20b^2\beta^2 - 24b^4\beta^4 - 6m_1 - 18b^2\beta^2m_1 \\ &+ (48\beta^2 + 8b^2\beta^4 - 48b^4\beta^6 - 24\beta^2m_1 \\ &- 12b^2\beta^4m_1 - 3\beta^2m_1^2 - 9b^2\beta^4m_1^2)s, \end{aligned}$$

$$\begin{aligned} n(C_b) &= 2b^4\beta^2 + 2b^2(m_1 - 1) \\ &+ (4b^4\beta^4 - 8b^2\beta^2 + 8b^2\beta^2m_1 + m_1^2)s, \end{aligned}$$

$$\begin{aligned} d(A_b) &= 2d(B_b) = 12\beta^2d(C_b) \\ &= 24\beta^2[1 + 3b^2\beta^2 + (4\beta^2 + 6b^2\beta^4)s], \end{aligned}$$

$$\omega_b(1) = 4/9, \quad \omega_b(0) = 0. \quad (\text{A10})$$

Diagram (c).

$$\begin{aligned} A_c(s) &= 64b^4\beta^2 + 12b^6\beta^4 - 60b^4\beta^2m_1 + 21b^2m_1^2 \\ &+ 36b^4\beta^2m_1^2 + (320b^2\beta^2 + 96b^4\beta^4 \\ &- 192b^2\beta^2m_1 + 24m_1^2 + 72b^2\beta^2m_1^2)s, \end{aligned}$$

$$\begin{aligned} B_c(s) &= 256b^4\beta^2 + 12b^6\beta^4 - 132b^4\beta^2m_1 \\ &+ 21b^2m_1^2 + 36b^4\beta^2m_1^2 + (320b^2\beta^2 + 96b^4\beta^4 \\ &- 192b^2\beta^2m_1 + 24m_1^2 + 72b^2\beta^2m_1^2)s, \end{aligned}$$

$$\begin{aligned} C_c(s) &= -32b^4\beta^2 + 4b^6\beta^4 + 32b^4\beta^2m_1 + 7b^2m_1^2 \\ &+ [32b^4\beta^4 + 64b^2\beta^2(m_1 - 1) + 8m_1^2]s, \end{aligned}$$

$$\eta_c(s) = 24\sqrt{3}b^3/\pi^3,$$

$$\begin{aligned} d(A_c) &= d(A_b) = 6d(A_c) \\ &= 48\beta^2[7b^2 + 6b^4\beta^2 + (8 + 12b^2\beta^2)s], \end{aligned}$$

$$d(\eta_c) = [7b^2 + 6b^4\beta^2 + (8 + 12b^2\beta^2)s]^{3/2},$$

$$\omega_c(1) = 4/9, \quad \omega_c(0) = 0. \quad (\text{A11})$$

Diagram (d).

$$\begin{aligned} n(A_d) &= 80b^2 + 160b^4\beta^2 + 12b^6\beta^4 - 72b^2m_1 - 132b^4\beta^2m_1 \\ &+ 21b^2m_1^2 + 36b^4\beta^2m_1^2 + (320b^2\beta^2 + 96b^4\beta^4 \\ &- 192b^2\beta^2m_1 + 24m_1^2 + 72b^2\beta^2m_1^2)s, \end{aligned}$$

$$n(B_d) = n(A_d),$$

$$n(C_d) = -16b^2 - 32b^4\beta^2 + 4b^6\beta^4 + 8b^2m_1 \\ + 20b^4\beta^2m_1 + b^2m_1^2 \\ + (-64b^2\beta^2 + 32b^4\beta^4 + 64b^2\beta^2m_1 + 8m_1^2)s,$$

$$n(\eta_d) = -24\sqrt{3}b^3\beta^3/\pi^3,$$

$$d(A_d) = d(B_d) = 6d(C_d) \\ = 48(2 + 3b^2\beta^2)(1 + 2b^2\beta^2 + 4\beta^2s),$$

$$d(\eta_d) = [(2 + 3b^2\beta^2)(1 + 2b^2\beta^2 + 4\beta^2s)]^{3/2},$$

$$\omega_d(1) = -4/9, \quad \omega_d(0) = 0. \quad (\text{A12})$$

2. Spin orbit

$$V_D^{SO}(\vec{p}, \vec{p}') = i(\vec{p} \times \vec{p}') \cdot \vec{S}_N w_D^{SO}(\vec{p}, \vec{p}'), \quad (\text{A13})$$

with S_N the spin operator of the nucleon and

$$w_D^{SO}(\vec{p}, \vec{p}') = 4\pi\alpha_s\omega_D(I) \int_0^\infty ds \eta_D(s) \\ \times \exp[-A_D(s)p^2 - B_D(s)p'^2 + C_D(s)\vec{p} \cdot \vec{p}'], \quad (\text{A14})$$

where the functions A, B, C are the same as in the Coulomb potential, and the η_D 's are given by

$$\eta_a(s) = \frac{3\sqrt{3}b^5(8-3m_1)}{32\pi^3[3b^2 + (2+3b^2\beta^2)s]^{5/2}},$$

$$\eta_b(s) = -\frac{3}{8\pi^3} \sqrt{\frac{3}{2}} \frac{b^3\beta^3(m_1-2)(4b^2\beta^2+m_1)}{[1+3b^2\beta^2 + (4\beta^2+6b^2\beta^4)s]^{5/2}},$$

$$\eta_c(s) = -\frac{3\sqrt{3}b^5(m_1-4-2b^2\beta^2)}{\pi^3[7b^2+6b^4\beta^2 + (8+12b^2\beta^2)s]^{5/2}},$$

$$\eta_d^{(1)}(s) = -\frac{3\sqrt{3}b^4\beta^4(1+2b^2\beta^2)(2b^2\beta^2+m_1)(28b^4\beta^4+m_1^2+8b^2\beta^2+8b^2\beta^2m_1)}{2\pi^3(2+3b^2\beta^2)(1+4b^2\beta^2+3b^4\beta^4)^{3/2}(4+2b^2\beta^2-m_1)[1+2\beta^2(b^2+2s)]^{5/2}},$$

$$\eta_d^{(2)}(s) = \frac{3\sqrt{3}b^3\beta^3(2b^2\beta^2+m_1)(28b^4\beta^4+m_1^2+8b^2\beta^2+8b^2\beta^2m_1)}{2\pi^3(2+3b^2\beta^2)^{5/2}(4b^2\beta^2+m_1)[1+2\beta^2(b^2+2s)]^{5/2}},$$

$$\omega_a(0) = -\frac{4}{3} \frac{1}{m^2}, \quad \omega_a(1) = +\frac{8}{9} \frac{1}{m^2},$$

$$\omega_b(0) = +\frac{2}{9} \left(\frac{1}{m_s^2} - \frac{1}{m^2} \right), \quad \omega_b(1) = -\frac{4}{27} \left(\frac{1}{m_s^2} - \frac{1}{m^2} \right),$$

$$\omega_c(0) = +\frac{1}{3} \frac{1}{m^2}, \quad \omega_c(1) = -\frac{1}{9} \frac{1}{m^2},$$

$$\omega_d^{(1)}(0) = -\frac{1}{9} \left(\frac{1}{m^2} + \frac{4}{mm_s} \right),$$

$$\omega_d^{(1)}(1) = -\frac{1}{27} \left(\frac{1}{m^2} + \frac{8}{mm_s} \right), \quad (\text{A15})$$

$$\omega_d^{(2)}(0) = +\frac{2}{9} \left(\frac{1}{m_s^2} + \frac{1}{mm_s} \right),$$

$$\omega_d^{(2)}(1) = -\frac{2}{27} \left(\frac{2}{m_s^2} - \frac{1}{mm_s} \right). \quad (\text{A16})$$

3. Contact spin-spin

$$V_D^{SS}(\vec{p}, \vec{p}') = \kappa_{ss}\omega_D(I) \eta_D \exp[-A_D p^2 - B_D p'^2 + C_D \vec{p} \cdot \vec{p}']. \quad (\text{A17})$$

Diagram (a).

$$A_a = \frac{2(1+\rho)^2 + 3g}{12\alpha^2(1+\rho)^2}, \quad B_a = A_a, \quad C_a = \frac{2(1+\rho)^2 + 3g}{6\alpha^2(1+\rho)^2},$$

$$\eta_a = 1, \quad \omega_a(1) = 1/3, \quad \omega_a(0) = 0,$$

$$\omega_a(1) = 1/3, \quad \omega_a(0) = 0.$$

Diagram (b).

$$A_b = \frac{(5g+3)\rho^2 + (-2g+6)\rho + (2g+3)}{6\alpha^2(g+3)(1+\rho)^2},$$

$$B_b = \frac{(5g+3)\rho^2 + (10g+6)\rho + (5g+3)}{6\alpha^2(g+3)(1+\rho)^2},$$

$$C_b = \frac{(1-g)\rho^2 + 2\rho + (g+1)}{\alpha^2(g+3)(1+\rho)^2},$$

$$\eta_b = \rho \left(\frac{6}{g+3} \right)^{3/2}, \quad \omega_b(1) = 1/3, \quad \omega_b(0) = 0.$$

Diagram (c).

$$A_c = \frac{(16g+3)\rho^2 + (2g+6)\rho + (21g^2+22g+3)}{12\alpha^2(7g+6)(1+\rho)^2},$$

$$B_c = \frac{(64g+3)\rho^2 + (62g+6)\rho + (21g^2+34g+3)}{12\alpha^2(7g+6)(1+\rho)^2},$$

$$C_c = \frac{(1-8g)\rho^2 + 2\rho + (7g^2+8g+1)}{2\alpha^2(7g+6)(1+\rho)^2},$$

$$\eta_c = \left(\frac{12g}{7g+6}\right)^{3/2}, \quad \omega_c(1) = 1/18, \quad \omega_c(0) = 1/6.$$

Diagram (d).

$$A_d = \frac{(20g^2+40g+3)\rho^2 + (4g^2+14g+6)\rho + (5g^2+10g+3)}{12\alpha^2(2g+3)(g+2)(1+\rho)^2},$$

$$B_d = A_d,$$

$$C_d = \frac{(1-4g^2-8g)\rho^2 + (2-4g^2-6g)\rho + (g^2+2g+1)}{2\alpha^2(2g+3)(g+2)(1+\rho)^2},$$

$$\eta_d = \rho \left[\frac{12g}{(2g+3)(g+2)} \right]^{3/2}, \quad \omega_d(1) = 1/18, \quad \omega_d(0) = 1/6.$$

4. Contact spin-independent

$$V_D^{\text{Con-SI}}(\vec{p}, \vec{p}') = 4\pi\alpha_s \omega_D(I) \exp[-A_D p^2 - B_D p'^2 + C_D \vec{p} \cdot \vec{p}'], \quad (\text{A18})$$

where the functions A_D, \dots are the same as for the spin-spin interaction and

$$\omega_a(1) = +\frac{1}{9} \frac{1}{m^2}, \quad \omega_a(0) = 0,$$

$$\omega_b(1) = -\frac{1}{18} \frac{1}{m^2} (1 + \rho^2), \quad \omega_a(0) = 0,$$

$$\omega_c(1) = -\frac{1}{9} \frac{1}{m^2}, \quad \omega_a(0) = 0,$$

$$\omega_c(1) = -\frac{1}{18} \frac{1}{m^2} (1 + \rho^2), \quad \omega_a(0) = 0. \quad (\text{A19})$$

5. Confinement

The confining interaction is taken to be a linearly rising potential, which in momentum space is given as

$$V_{conf} = \frac{6\pi\sigma}{q^4}, \quad (\text{A20})$$

where σ is the string tension. Equation (A20) includes a color factor of 3/4. The effective KN interaction is given by

$$V_D^{\text{Conf}}(\vec{p}, \vec{p}') = 6\pi\sigma \omega_D(I) \int_0^\infty du \int_u^\infty ds \eta_D(s) \times \exp[-A_D(s)p^2 - B_D(s)p'^2 + C_D(s)\vec{p} \cdot \vec{p}'], \quad (\text{A21})$$

where the functions $A_D(s), B_D(s), C_D(s)$, and η_D are the same as for the Coulomb term, and the $\omega_D(I)$'s are given by

$$\omega_a(1) = +4/9, \quad \omega_a(0) = 0,$$

$$\omega_b(1) = -4/9, \quad \omega_b(0) = 0,$$

$$\omega_c(1) = -4/9, \quad \omega_a(0) = 0,$$

$$\omega_d(1) = +4/9, \quad \omega_d(0) = 0. \quad (\text{A22})$$

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