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Study of the $\bar{D}N$ Interaction in a QCD Coulomb Gauge Quark Model

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Abstract. We study the $\bar{D}N$ interaction at low energies with a quark model inspired in the QCD Hamiltonian in Coulomb gauge. The model Hamiltonian incorporates a confining Coulomb potential extracted from a self-consistent quasiparticle method for the gluon degrees of freedom, and transverse-gluon hyperfine interaction consistent with a finite gluon propagator in the infrared. Initially a constituent-quark mass function is obtained by solving a gap equation and baryon and meson bound-states are obtained in Fock space using a variational calculation. Next, having obtained the constituent-quark masses and the hadron waves functions, an effective meson-nucleon interaction is derived from a quark-interchange mechanism. This leads to a short range meson-baryon interaction and to describe long-distance physics vector- and scalar-meson exchanges described by effective Lagrangians are incorporated. The derived effective $\bar{D}N$ potential is used in a Lippmann-Schwinger equation to obtain phase shifts. The results are compared with a recent similar calculation using the nonrelativistic quark model.

1 Introduction

The study of the DN interaction is of interest in different contexts, in particular in problems of chiral symmetry restoration in hadronic matter. The chiral properties of the light u and d quarks in D mesons are very sensitive to temperature and baryon density and one can expect modifications in the properties of these mesons in medium. For example, in-medium spectral functions of D mesons [1] require a detailed model for DN interaction and, as emphasized in recent publications [2,3], the understanding of this interaction in free space is a prerequisite for constraining such calculations.

In Refs. [2,3] the free-space $\bar{D}N$ interaction ($\bar{D}$ stands for $\bar{D}^0$ and D^-) was studied in the context of a hybrid model, where the short-distance part of the interaction was described by quark-gluon interchange in a nonrelativistic constituent quark model, and the long-distance part was described by meson-baryon exchange. $\bar{D}N$ reactions are easier to treat within a quark model than D^+N reactions, because in $\bar{D}N$ there are no direct quark-antiquark annihilation contributions to the scattering matrix, only quark-gluon interchange contribute. However, although quark-gluon interchange processes are efficiently described within a nonrelativistic quark model, such a model is clearly very limited for studying in-medium chiral effects. In a typical nonrelativistic model there is no physical link between the masses of the constituent quarks, which are supposed to be generated by the phenomenon of dynamical chiral symmetry breaking (DχSB), and the microscopic quark-gluon interaction, which is responsible for the dynamical breaking of the symmetry; the masses and the interactions are specified independently in the model. For addressing in medium chiral effects one needs a model that is closer to QCD, but it should be still simple enough to allow perform calculations without resorting to extensive numerical simulations.

A model that provides such a workable framework is based on the QCD Hamiltonian in Coulomb gauge [4]. The model is based on a field theoretic Hamiltonian that confines color and realizes DχSB [5]. Color confinement and DχSB are supposedly the most prominent nonperturbative effects in QCD, and are thought to be responsible for the properties of the structure and interactions of hadrons at low energies. The model allows to construct an approximation scheme for obtaining hadron wave functions and effective hadron-hadron potentials that can be iterated in a Lippmann–Schwinger equation to calculate scattering observables [6]. In the present communication we sketch the derivation of such a model, and discuss its use to derive the short-distance part of the $\bar{D}N$ interaction in free space. In order to compare with the results obtained in Ref. [2], we add long-distance contributions from one-meson exchanges.

The paper is organized as follows. In the next Section we introduce the model Hamiltonian, discuss DχSB, and describe an approximation scheme to obtain hadron wave functions. Having obtained the constituent-quark masses with DχSB and the hadron waves functions, in Section 3 we derive an effective meson-baryon interaction through a quark-interchange mechanism. In Section 4 we add to the quark-model interaction long-distance meson exchanges. Numerical results for the constituent-quark mass function, hadron wave functions and $\bar{D}N$ phases shifts are presented in Section 5. Our Conclusions and Perspectives are presented in Section 6. In Appendix A we present the effec-

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tive meson-baryon Lagrangian densities used in the present work.

2 Dynamical Chiral Symmetry Breaking and a Constituent Quark Model

Most of the nonrelativistic constituent quark models share the common feature that hadrons are bound states of massive constituent quarks and antiquarks confined by a prescribed potential. In addition, residual interactions derived from perturbative one-gluon exchange (OGE) are added to describe spin-dependent hadron mass splittings. In this Section we show how such a constituent quark model can be derived [6] from an underlying Hamiltonian that confines color and realizes DχSB [5]. We also discuss a calculation scheme for building hadron states in a Fock space, in close analogy to the nonrelativistic quark model. Given the microscopic interactions and the Fock-space hadronic states, low-energy effective hadronic interactions can be derived using a resonating group method (RGM) - for a review and a list of references on the RGM and other methods, see Ref. [7].

The starting point is the specification of a microscopic Hamiltonian. Here we take an Hamiltonian inspired in the Coulomb gauge QCD Hamiltonian, specifically [5]

$$\begin{aligned} H = & \int d\mathbf{x}\, \Psi^\dagger(\mathbf{x})[-i\alpha\cdot\nabla + \beta\, m_0]\Psi(\mathbf{x}) \\ & - \frac{1}{2}\int d\mathbf{x}\, d\mathbf{y}\, \rho^a(\mathbf{x})\, V_C(|\mathbf{x}-\mathbf{y}|)\, \rho^a(\mathbf{y}) \\ & + \frac{1}{2}\int d\mathbf{x}\, d\mathbf{y}\, J^a_i(\mathbf{x})\, V_{ij}(|\mathbf{x}-\mathbf{y}|)\, J^a_j(\mathbf{y}), \end{aligned} \tag{1}$$

where $\Psi(\mathbf{x})$ is the quark field operator (with color and flavor indices suppressed), m_0 is the current-quark mass matrix, and $\rho^a(\mathbf{x})$ and $J^a_i(\mathbf{x})$ are the charge and current densities given by

$$\rho^a(\mathbf{x}) = \Psi^\dagger(\mathbf{x})\, T^a\, \Psi(\mathbf{x}), \quad J^a_i(\mathbf{x}) = \Psi^\dagger(\mathbf{x})\, T^a\, \alpha_i \Psi(\mathbf{x}) \tag{2}$$

with $T^a = \lambda^a/2, a = 1, \cdots, 8$, where λ^a are the color SU(3) Gell-Mann matrices. In Eq. (1), $V_C(|\mathbf{x}-\mathbf{y}|)$ is the "Coulomb potential", and $V_{ij}(|\mathbf{x}-\mathbf{y}|)$ is a transverse-gluon hyperfine potential

$$V_{ij}(|\mathbf{x}-\mathbf{y}|) = \left(\delta_{ij} - \frac{\nabla_i\nabla_j}{\nabla_j}\right) V_T(|\mathbf{x}-\mathbf{y}|). \tag{3}$$

Given V_C and V_T, the microscopic Hamiltonian is fully specified.

DχSB can be discussed in a Bogoliubov-Valatin approach. Initially, the quark field operator is expanded in a basis of plane-wave spinors u and v as

$$\Psi(\mathbf{x}) = \int \frac{d\mathbf{x}}{(2\pi)^{3/2}} \sum_{s=\pm 1/2} [u_s(\mathbf{k})q_s(\mathbf{k}) + v_s(\mathbf{k})\bar{q}^\dagger_s(-\mathbf{k})]\, e^{i\mathbf{k}\cdot\mathbf{r}}, \tag{4}$$

where

$$u_s(\mathbf{k}) = \sqrt{\frac{E_k + M_k}{2E_k}} \begin{pmatrix} 1 \\ \frac{\sigma\cdot\mathbf{k}}{E_k+M_k} \end{pmatrix} \chi_s, \tag{5}$$

and

$$v_s(\mathbf{k}) = \sqrt{\frac{E_k + M_k}{2E_k}} \begin{pmatrix} -\frac{\sigma\cdot\mathbf{k}}{E_k+M_k} \\ 1 \end{pmatrix} \chi^c_s, \tag{6}$$

with

$$E_k = \sqrt{k^2 + M^2_k}, \tag{7}$$

where M_k is the constituent-quark mass function, χ_s is a Pauli spinor, and $\chi^c_s = -i\sigma^2\chi^*_s$. The operators $q^\dagger_s(\mathbf{k})$, $q^\dagger_s(\mathbf{k})$, $\bar{q}_s(\mathbf{k})$, and $\bar{q}_s(\mathbf{k})$ are creation and annihilation operators of *constituent quarks*, i.e.

$$q_s(\mathbf{k})|\Omega\rangle = 0, \qquad \bar{q}_s(\mathbf{k})|\Omega\rangle = 0, \tag{8}$$

where $|\Omega\rangle$ is the chirally broken vacuum state. That is, when $m_0 = 0$ the Hamiltonian in Eq. (1) is chirally symmetric but the vacuum state breaks this symmetry,

$$\langle\Omega|\bar{\Psi}\Psi|\Omega\rangle \neq 0. \tag{9}$$

The momentum-dependent constituent-quark mass function M_k is determined as follows. First, the field operator Ψ of Eq. (4) and the corresponding $\Psi^\dagger$ are substituted in the Hamiltonian of Eq. (1), and Wick's theorem is used to bring the Hamiltonian to normal order with respect to the constituent quark operators. This leads to a Hamiltonian of the following structure

$$H = \mathcal{E} + H_2 + H_4, \tag{10}$$

where $\mathcal{E}$ is a c-number function, the vacuum energy, and H_2 and H_4 contain respectively normal-ordered bilinear and quadrilinear terms in the quark field operators Ψ. Then, one requires that H_2 is diagonal in the constituent quark and antiquark creation and annihilation operators, which leads to the *gap equation* for the mass function M_k,

$$\begin{aligned} M_k = m_0 + \frac{2}{3}\int \frac{d\mathbf{q}}{(2\pi)^3 E_q} \Big[& f_1(\mathbf{k},\mathbf{q})\, V_C(|\mathbf{k}-\mathbf{q}|) \\ & + g_1(\mathbf{k},\mathbf{q})\, V_T(|\mathbf{k}-\mathbf{q}|)\Big], \end{aligned} \tag{11}$$

with

$$f_1(\mathbf{k},\mathbf{q}) = M_q - M_k\frac{q}{k}\,\hat{k}\cdot\hat{q}, \tag{12}$$

$$g_1(\mathbf{k},\mathbf{q}) = 2M_q + 2M_k\frac{q}{k}\frac{(\mathbf{k}\cdot\mathbf{q}-k^2)(\mathbf{k}\cdot q - q^2)}{kq|\mathbf{k}-\mathbf{q}|^2}. \tag{13}$$

We note that the gap equation can also be obtained by minimizing the vacuum energy $\mathcal{E}$ with respect to the constituent-quark mass-function.

Once the gap equation is satisfied, H_2 is diagonal in the quark and antiquark operators

$$H_2 = \int d\mathbf{k}\,\varepsilon_k \sum_{s=\pm 1/2}\Big[\, q^\dagger_s(\mathbf{k})q_s(\mathbf{k}) + \bar{q}^\dagger_s(\mathbf{k})\bar{q}_s(\mathbf{k})\Big], \tag{14}$$

where the constituent-quark single-particle energy is given by

$$\varepsilon_k = \frac{k^2 + m_0 M_k}{E_k} + \frac{2}{3}\int \frac{d\mathbf{q}}{(2\pi)^3 E_q}\Big[f_2(\mathbf{k},\mathbf{q})\, V_C(|\mathbf{k}-\mathbf{q}|) + g_2(\mathbf{k},\mathbf{q})\, V_T(|\mathbf{k}-\mathbf{q}|)\Big], \quad (15)$$

with

$$f_2(\mathbf{k},\mathbf{q}) = M_k M_q + \hat{k}\cdot\hat{q}, \quad (16)$$

$$g_2(\mathbf{k},\mathbf{q}) = 2M_k M_q + 2\frac{(\mathbf{k}\cdot\mathbf{q}-k^2)(\mathbf{k}\cdot q - q^2)}{|\mathbf{k}-\mathbf{q}|^2}. \quad (17)$$

H_4 is simply the normal-ordered form of the last two terms of the Hamiltonian in Eq. (1), i.e.

$$H_4 = \frac{1}{2}\int d\mathbf{x}\, d\mathbf{y}\, V_\Gamma(\mathbf{x}-\mathbf{y}) \times :\Big[\Psi^\dagger(\mathbf{x})\,\Gamma\, T^a\, \Psi(\mathbf{x})\Big]\Big[\Psi^\dagger(\mathbf{y})\,\Gamma\, T^a\, \Psi(\mathbf{y})\Big]:, \quad (18)$$

where V_Γ is $-V_C$ for $\Gamma = 1$ and V_{ij} for $\Gamma = \alpha^i$.

Given the quark mass function, the hadron bound states can be determined in a variety of ways. For our purposes it is more convenient to work in Fock space. The one-baryon state in a BCS approximation can be written as [8]

$$|a\rangle = B_a^\dagger|\Omega\rangle = \frac{1}{\sqrt{3!}}\,\psi_a^{\mu_1\mu_2\mu_3}\, q^\dagger_{\mu_1} q^\dagger_{\mu_2} q^\dagger_{\mu_3}|\Omega\rangle, \quad (19)$$

with $\psi_a^{\mu_1\mu_2\mu_3}$ being the Fock-space amplitude, with the indices a and μ_1, μ_2, μ_3 denoting the quantum numbers of the baryon (orbital, spin, flavor) and of the quarks (orbital, color, spin, flavor), respectively. The one-meson state is defined analogously as

$$|a\rangle = M_a^\dagger|\Omega\rangle = \phi_a^{\mu\nu} q^\dagger_\mu \overline{q}^\dagger_\nu|\Omega\rangle, \quad (20)$$

where $\phi_a^{\mu\nu}$ is the corresponding Fock-space amplitude, with a and μ, ν representing the quantum numbers of the mesons and quarks and antiquarks, respectively. Given the hadron Fock-space states in the BCS approximation, the Fock amplitudes ψ and ϕ can be obtained by solving a Schrödinger-like equation. However, in order to simplify the derivation, a variational calculation is used [9,8]. This amounts to making an ansatz for the amplitudes and minimizing the expectation values of the Hamiltonian in the one-baryon and one-meson states with respect to variational parameters.

Specifically, we work in momentum space and write for the momentum-dependent part of ψ and ϕ the following Gaussian forms with variational parameters α and β

$$\psi_{\mathbf{P}}(\mathbf{k}_1,\mathbf{k}_2,\mathbf{k}_3) = \delta(\mathbf{P}-\mathbf{k}_1-\mathbf{k}_2-\mathbf{k}_3) \times \left(\frac{3}{\pi^2\alpha^4}\right)^{3/4} e^{-\sum_{i=1}^{3}(\mathbf{k}_i-\mathbf{P}/3)/2\alpha^2}, \quad (21)$$

and

$$\phi_{\mathbf{P}}(\mathbf{k}_1,\mathbf{k}_2) = \delta(\mathbf{P}-\mathbf{k}_1-\mathbf{k}_2) \times \left(\frac{1}{\pi\beta^2}\right)^{3/2} e^{-(M_1\mathbf{k}_1-M_2\mathbf{k}_2)^2/8\beta^2}, \quad (22)$$

where $\mathbf{P}$ is the center-of-mass momentum of the hadrons,

$$M_1 = \frac{2M_{\bar q}}{M_q + M_{\bar q}}, \quad (23)$$

and

$$M_2 = \frac{2M_q}{M_q + M_{\bar q}}, \quad (24)$$

with M_q and $M_{\bar q}$ being constituent-quark masses. Their values will be specified in the following. As said, the variational parameters α and β are determined minimizing the meson and baryon masses,

$$M_a = \left.\frac{\langle a|\,(H_2 + H_4)\,|a\rangle}{\langle a|a\rangle}\right|_{\mathbf{P}_a=0}. \quad (25)$$

In principle, given M_k one can calculate the expectation value in Eq. (25). However, the calculation of the expectation value is still too complicated, since it involves multi-dimensional integrals of products of the spinors u and v coming from H_4. Fortunately, only the low-momentum part of M_k in the expressions for u and v is required because the rapid falloff of wavefunctions ϕ and ψ cut off the integrand in Eq. (25) at large values of the momenta. This means that one can make an approximation for M_k as

$$M_k \simeq M - M_1 k - M_2 k^2, \quad (26)$$

where $M = M_{k=0}$. The constituent-quark masses M_q and $M_{\bar q}$ in Eqs. (23) and (24) are then simply the zero-momentum component of the corresponding mass functions. Replacing the expansion of Eq. (26) in the expressions for the spinors u and v in Eqs. (5) and (6), and retaining terms up to $O(k^2/M^2)$, one obtains for u and v the following expressions

$$u_s(\mathbf{k}) \simeq \begin{pmatrix} 1 - \frac{k^2}{8M^2} \\ \frac{\sigma\cdot\mathbf{k}}{2M} \end{pmatrix} \chi_s, \quad (27)$$

and

$$v_s(\mathbf{k}) \simeq \begin{pmatrix} -\frac{\sigma\cdot\mathbf{k}}{2M} \\ 1 - \frac{k^2}{8M^2} \end{pmatrix} \chi^c_s. \quad (28)$$

Using these expressions to evaluate the expectation value in Eq. (25), one obtains the following nonlinear equations for the variational parameters α and β (for $M_u = M_d$)

$$\alpha^4 = \frac{M_u}{3}\int \frac{d\mathbf{q}}{(2\pi)^3}\, q^2\,\langle V(\mathbf{q})\rangle_N\, e^{-q^2/2\alpha^2} \quad (29)$$

and

$$\beta^4 = \frac{2}{9}\left(\frac{M_u M_c}{M_u + M_c}\right)\int \frac{d\mathbf{q}}{(2\pi)^3}\, q^2\,\langle V(\mathbf{q})\rangle_{\bar D}\, e^{-q^2/4\beta^2}, \quad (30)$$

where we keep only the dominant contributions of the microscopic quark-gluon interaction, namely the central component of the Coulomb interaction and the spin-spin component of the transverse-gluon interaction,

$$V(\mathbf{q}) = V_C(\mathbf{q}) + \frac{2}{3}\frac{q^2}{M_1 M_2}\mathbf{S}_1\cdot\mathbf{S}_2\, V_T(\mathbf{q}). \quad (31)$$

For the nucleon, we have

$$\langle V(\mathbf{q})\rangle_N = V_C(q) + \frac{2}{3}\frac{q^2}{M_u^2}\left(-\frac{1}{4}\right)V_T(q), \qquad (32)$$

and for the $\bar{D}$ meson

$$\langle V(\mathbf{q})\rangle_{\bar{D}} = V_C(q) + \frac{2}{3}\frac{q^2}{M_u M_c}\left(-\frac{3}{4}\right)V_T(q). \qquad (33)$$

Note that in obtaining these expressions, M_q and E_q in the integrals of the functions $f_2(\mathbf{k},\mathbf{q})$ and $g_2(\mathbf{k},\mathbf{q})$ in Eqs. (16) and (17) are expanded consistently up to second order in the momenta.

3 Quark-Model Meson-Baryon Interaction

Given the microscopic Hamiltonian and the hadronic states in Fock space an effective hadron-hadron interaction can be derived through a RGM-type formalism [7]. The part of the Hamiltonian relevant for the $\bar{D}N$ interaction can be written in a compact notation as

$$H_{q\bar{q}} = T(\mu)\, q^\dagger_\mu q_\mu + T(\nu)\, \bar{q}^\dagger_\nu \bar{q}_\nu + \frac{1}{2} V_{qq}(\mu\nu;\sigma\rho)\, q^\dagger_\mu q^\dagger_\nu q_\rho q_\sigma + \frac{1}{2} V_{\bar{q}\bar{q}}(\mu\nu;\sigma\rho)\, \bar{q}^\dagger_\mu \bar{q}^\dagger_\nu \bar{q}_\rho \bar{q}_\sigma + V_{q\bar{q}}(\mu\nu;\sigma\rho)\, q^\dagger_\mu \bar{q}^\dagger_\nu \bar{q}_\rho q_\sigma, \qquad (34)$$

with the indices (μ,ν,ρ,σ) denoting as before the quantum numbers of quarks and antiquarks. The V_{qq}, $V_{q\bar{q}}$, and $V_{\bar{q}\bar{q}}$ come from the contraction of the spinors in Eq. (18) and are proportional to V_C and V_T. Their explicit expressions will be given elsewhere.

For a $H_{q\bar{q}}$ in this form, the effective meson-baryon potential for the process $a + b \to c + d$ can be written as [7]

$$\begin{aligned} V(ab,cd) = &-3\,\phi_c^{*\mu\nu_1}\,\psi_d^{*\nu\mu_2\mu_3}\,V_{qq}(\mu\nu;\sigma\rho)\,\phi_a^{\rho\nu_1}\,\psi_b^{\sigma\mu_2\mu_3} \\ &-3\,\phi_c^{*\sigma\rho}\,\psi_d^{*\mu_1\mu_2\mu_3}\,V_{q\bar{q}}(\mu\nu;\sigma\rho)\,\phi_a^{\mu_1\rho}\,\psi_b^{\mu\mu_2\mu_3} \\ &-6\,\phi_c^{*\mu_1\nu_1}\,\psi_d^{*\nu\mu\mu_3}\,V_{qq}(\mu\nu;\sigma\rho)\,\phi_a^{\rho\nu_1}\,\psi_b^{\mu_1\sigma\mu_3} \\ &-6\,\phi_c^{*\nu_1\nu}\,\psi_d^{*\nu_1\mu\mu_3}\,V_{q\bar{q}}(\mu\nu;\sigma\rho)\,\phi_a^{\nu_1\rho}\,\psi_b^{\mu_1\sigma\mu_3}, \end{aligned} \qquad (35)$$

where $\phi_a^{\mu\nu}$, $\psi_b^{\mu\nu\sigma}$, $\cdots$ are the meson and baryon Fock-space amplitudes discussed in the previous Section. As previously, there is a sum or and integral over repeated quark indices.

The expression in Eq. (35) for the effective meson-baryon potential is completely general, valid for any meson-baryon process for which the baryon and meson Fock-space states are given as in Eqs. (19) and (20), and for a generic H_4 in form of Eq. (18). In this work we are interested in the interaction of $\bar{D}$ mesons and nucleons, more specifically in the elastic scattering channels $p\bar{D}^0 \to p\bar{D}^0$, $nD^- \to nD^-$, $pD^- \to pD^-$, $n\bar{D}^0 \to n\bar{D}^0$, and in the charge-exchange channel $pD^- \to n\bar{D}^0$. In order to avoid multidimensional numerical integrations, we use the low momentum approximation for the constituent quark mass function M_k and the Gaussian wave functions with the variational parameters α and β determined by minimizing the nucleon and $\bar{D}$ mesons, as discussed in the previous Section. In addition, use the same microscopic quark interaction as used in the bound-state calculation, namely, the central component of the Coulomb interaction and the spin-spin component of the transverse-gluon interaction. With these ingredients, the effective $\bar{D}N$ interaction can be written as a sum of four contributions (each of these can be represented pictorially by quark-interchange graphs as in Fig. 1)

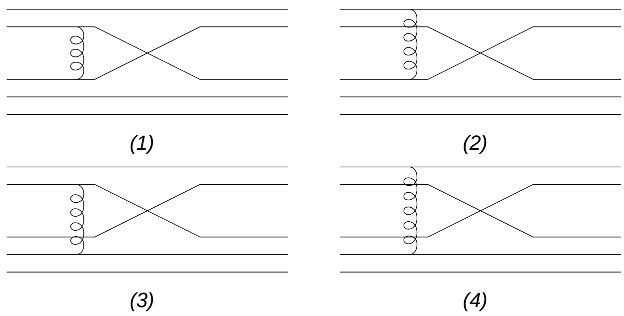

Fig. 1. Pictorial representation of the quark-interchange processes that contribute to a meson-baryon interaction. The curly lines represent the gluon interactions V_C and V_T.

$$V_{\bar{D}N}(\mathbf{p},\mathbf{p}') = \frac{1}{2}\sum_{i=1}^{4}\omega_i\,[\,V_i(\mathbf{p},\mathbf{p}') + V_i(\mathbf{p}',\mathbf{p})\,], \qquad (36)$$

with

$$V_i(\mathbf{p},\mathbf{p}') = \left[\frac{3g}{(3+2g)\pi\alpha^2}\right] e^{-c_i p^2 - d_i p'^2 + e^i \mathbf{p}\cdot\mathbf{p}'} \times \int \frac{d\mathbf{q}}{(2\pi)^3}\, v(q)\, e^{-a_i q^2 + \mathbf{b}_i\cdot\mathbf{q}}, \qquad (37)$$

where the a_i, b_i, c_i, d_i, and e_i involve the hadron sizes α and β, and the quark masses M_u and M_c. Since we are using the same type of wave functions as in Ref. [2], their explicit expressions are the same as in that reference (with the a_i interchanged with the c_i, etc). In addition, in Eq. (37), $v(q)$ is $V_C(q)$ or $2q^2/(3M_1M_2)V_T(q)$. The ω_i come from the sum over quark color-spin-flavor indices and combinatorial factors, whose values are given in Table 1.

Table 1. Coefficients ω_i of color, spin and flavor due to central $1_i 1_j$ and spin-spin $S_i S_j$ interaction for the isospin $I = 0$ and $I = 1$ states.

Isospin	ω_1		ω_2		ω_3		ω_4	
	$1_i 1_j$	$S_i S_j$	$1_i 1_j$	$S_i S_j$	$1_i 1_j$	$S_i S_j$	$1_i 1_j$	$S_i S_j$
$I = 0$	0	0	0	0	0	-1/6	0	-1/6
$I = 1$	-4/9	-1/3	4/9	-1/3	4/9	-1/18	-4/9	-1/18

4 Meson-Exchange contributions to the $\bar{D}N$ Interaction

The quark-interchange processes are very short ranged. It is known that in other low-energy hadron-hadron reactions, long-range meson-exchange processes play an important role and one expects that they will equally be important for the present $\bar{D}N$ reaction. For example, in the case of the K^+N reaction, quark interchange accounts only for part of the experimental s-wave phase shifts and σ, ω and ρ meson exchanges are crucial for describing this wave as well as higher partial waves [11]. It is important also to add that meson-baryon dynamics alone [12, 13] seems not to be sufficient to describe the experimental K^+N data. With these facts in mind, Ref. [2] extended the approach of Ref. [11] for the K^+N system to the $\bar{D}N$ system. The conclusion of Ref. [2] was that the dominant contributions from meson-exchanges come from vector ρ and ω exchanges, and scalar contributions that we parametrize here in terms of a single σ-meson exchange – see Fig. 2. We note that the replacement of a correlated $\pi\pi$ exchange by a single σ-meson exchange is not a bad approximation for the $I = 0$ channel, but for $I = 1$ channel it provides only 50% of the total strength [2].

The interacting meson-meson and meson-baryon Lagrangian densities we use are given in Appendix A. The tree-level potentials derived from the Lagrangian densities lead to the following expressions, for the vector-meson exchanges ($v = \rho, \omega$)

$$V^v(\mathbf{p}', \mathbf{p}) = \frac{g_{NNv}\, g_{\bar{D}\bar{D}v}}{\sqrt{4\omega(p)\omega(p')}}\,(p' + p)_\mu\, \varDelta_v^{\mu\nu}(q) \times \left[A_\nu(p's', ps) + \left(\frac{\kappa_v}{2m_v}\right) B_\nu(p's', p, s)\right], \quad (38)$$

and for the σ exchange

$$V^\sigma(\mathbf{p}', \mathbf{p}) = \frac{g_{NN\sigma}\, g_{\bar{D}\bar{D}\sigma}}{\sqrt{4\omega(p)\omega(p')}}\, \varDelta_\sigma(q)\bar{u}(\mathbf{p}', s')u(\mathbf{p}, s), \quad (39)$$

where

$$\omega(q) = \sqrt{q^2 + m_D^2}, \quad (40)$$

and

$$A_\mu(p's', ps) = \bar{u}(\mathbf{p}', s')\gamma_\mu u(\mathbf{p}, s), \quad (41)$$

and

$$B_\mu(p's', p, s) = \bar{u}(\mathbf{p}', s')\, i\sigma_{\mu\nu}\, q^\nu\, u(\mathbf{p}, s). \quad (42)$$

Here, $\varDelta_v^{\mu\nu}(q)$ and $\varDelta_\sigma(q)$ are the vector-meson and scalar-meson propagators, and $u(\mathbf{p}, s)$ and $\bar{u}(\mathbf{p}, s)$ are the Dirac spinors of nucleons (not of quarks).

To avoid divergences in the Lippmann-Schwinger equation, we introduce as usual form factors at each vertex in Fig. 2. The form factors are of monopole type, given by

$$F_i(\mathbf{q}^2) = \left(\frac{\varLambda_i^2 - m_i^2}{\varLambda_i^2 + \mathbf{q}^2}\right) \quad (43)$$

where $\mathbf{q} = \mathbf{p}' - \mathbf{p}$, m_i is the boson mass in the propagator and $\varLambda_i$ is a cutoff mass.

The values of the coupling constants, masses and cutoff masses are discussed in the next section.

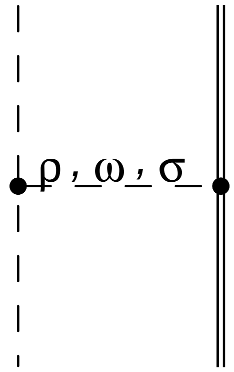

Fig. 2. One-meson exchanges contributions to $\bar{D}N$ interaction.

5 Numerical Results

The first step is the specification of the interactions V_C and V_T of the microscopic Hamiltonian. For simplicity, we take for V_C the interaction determined by Szczepaniak and Swanson in Ref. [14] using a self-consistent quasiparticle method for the gluon degrees of freedom in Coulomb gauge QCD. Namely, V_C can be fitted in momentum space to the full numerical solution as [14]

$$V_C(q) = V_l(q) + V_s(q), \quad (44)$$

with

$$V_l(q) = \frac{(3.50)^2}{q^2}\left(\frac{m_g}{q}\right)^{1.93}, \quad (45)$$

for $q \leq m_g$, and

$$V_s(q) = \frac{8.07/q^2}{\ln^{0.8}(q^2/m_g^2 + 1.41)\, \ln^{0.62}(q^2/m_g^2 + 0.82)}, \quad (46)$$

for $q > m_g$. The gluonic mass m_g is a renormalization scale which was fixed to $m_g = 600$ MeV by fitting the lattice heavy-quark potential. We could also have used a more recent fit to a full numerical solution of an improved calculation [15], but the final results would not be much different from the ones presented here. For the transverse-gluon hyperfine potential, we use the form

$$V_T(q) = -\frac{4\pi\alpha_T}{(q^2 + m^2)\ln^{1.42}(\tau + q^2/\mu_g^2)}, \quad (47)$$

where the parameters α_T, μ and τ are chosen to get a reasonable value for the quark condensate and a constituent quarks mass of $M_u = 300$ MeV. For that we use $\mu = m_g/2$, $\tau = 1.05$ and $\alpha_T = 0.7$. The transverse-gluon hyperfine interaction consistent with a finite gluon propagator in the infrared as seen in recent lattice QCD simulations in Coulomb gauge [16].

The numerical solution of the gap equation leads to the results shown in Fig. 3 for different values of the current mass m_0. Two important facts are to be noticed in this figure. First, for momenta $k \leq m_g = 600$ MeV, the mass function has very little momentum dependence (in a log scale). This is important in connection with the low-momentum

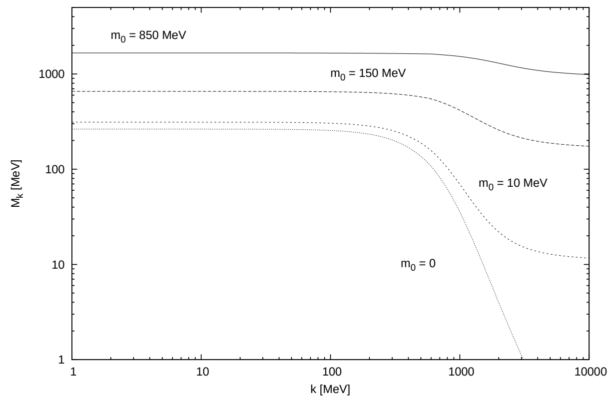

Fig. 3. Constituent quark mass M_k as a function of the momentum for different values of the current quark mass m_0.

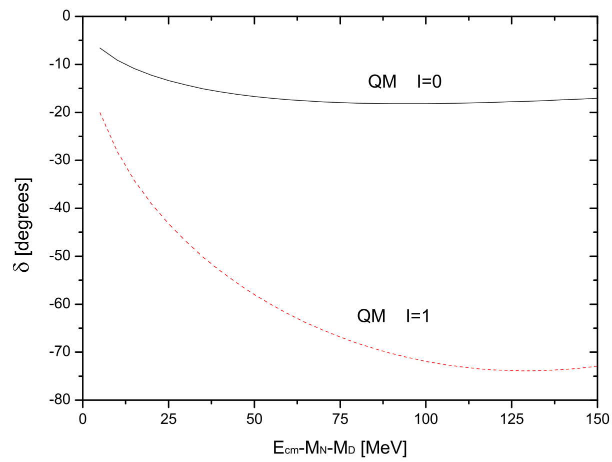

Fig. 4. s-wave phase-shifts for $I = 0$ and $I = 1$ channels for the Coulomb Gauge quark model (QM).

expansion of the mass function, Eq. (26), meaning that an approximation that uses only $M_{k=0}$ should be a good one if the wave functions cut off momentum components larger than m_g. Second, as m_0 increases, the effects of DχSB diminish, but even for a current quark mass as large as m_c = 850 MeV, there is still a significant dressing of the interactions.

The next step towards the determination of the $\bar{D}N$ interaction requires the determination of the values of the variational parameters of the nucleon (α) and the D meson (β). The numerical solution of Eqs. (29) and (30), for the same microscopic interactions V_C and V_T and the mass function determined above, leads to α = 485 MeV and β = 530 MeV. It is worth commenting that the standard quark model values, as employed in Ref. [3], are α = 383.5 MeV and β = 400 MeV. This means that our hadron sizes (the hadron r.m.s. radii are inversely proportional to α and β) are smaller than those of Ref. [3]. Since the sizes of the wave functions influence the degree of overlap of the colliding hadrons, the quark-interchange potentials we obtain here are expected to be different from those of Ref. [3]. For a recent discussion on the influence of hadron sizes on the quark-interchange mechanism, see Ref [17].

We are in position to discuss our numerical results for $\bar{D}N$ scattering. We present results for the s-wave phase-shifts for isospin $I = 0$ and $I = 1$. We solve the Lippmann-Schwinger equation for the potential given in Eqs. (36) and (37). We reiterate that we use the constituent-quark masses and size parameters α and β as determined above, and the same microscopic V_C and interactions used to determine these quantities. The results are presented Fig. 4 as function of the center-of-mass (c.m.) kinetic energy.

As always, quark-interchange leads to repulsive interactions. We note that the repulsion is larger in the present model than in the one-gluon-exchange model of Ref. [3]. One reason for this is the smaller size of the wave functions, which leads to a larger overlap between the wave functions of the hadrons. Another reason is that the interaction at small momenta in the present model is larger than the one-gluon-exchange and this also gives a larger value for the effective hadron-hadron interaction.

Next we discuss the role of meson exchange. The coupling constants are fixed used SU(4) symmetry [3], i.e. the coupling constants of the vector mesons to the charmed mesons are related to $g_{\pi\pi\rho}$ as

$$g_{\bar{D}\bar{D}\rho} = g_{\bar{D}\bar{D}\omega} = \frac{g_{\pi\pi\rho}}{2}. \tag{48}$$

We use the empirical value of $g_{\pi\pi\rho}$, namely $g_{\pi\pi\rho}$ = 6.0 For the vertices involving nucleons we use the following values $g_{NN\sigma}$ = 13.3, $g_{NN\rho}$ = 3.25, κ_ρ = 6.1, $g_{NN\omega}$ = 3.25. The values for mesons and baryon masses are m_p = 939 MeV, m_n = 937 MeV, m_ρ = 770 MeV, m_ω = 783.4 MeV and m_σ = 550 MeV. And finally, the numerical values for the cutoff masses are the same as in Ref. [3], namely $\Lambda_{NN\rho}$ = 1.6 GeV, $\Lambda_{\bar{D}\bar{D}\rho}$ = 0.917 GeV, $\Lambda_{NN\omega}$ = 2.750 GeV, $\Lambda_{\bar{D}\bar{D}\omega}$ = 0.843 GeV, $\Lambda_{NN\sigma}$ = 1.2 GeV, $\Lambda_{\bar{D}\bar{D}\sigma}$ = 1.7 GeV. Our results for the s-wave phase-shifts from meson-exchanges are shown in Figs. 5 and 6.

Clearly seen in Fig. 5 is the destructive interference of the ρ and ω contributions, a feature first noticed in Ref. [3]. Inclusion of the σ has little effect. However, for the $I = 1$ channel, shown in Fig. 6, both ρ and ω contributions are repulsive and there is constructive interference. The contribution of the one-σ exchange is very small, as in $I = 0$. Finally, in Fig. 7 we show the results considering both quark-interchange and meson-exchange. As expected, the effects are additive.

A natural question that might arise at this point is the potential for double counting when considering simultaneously quark-interchange and meson-exchange. In fact, only part of ω exchange (the 0-th component of the $\phi_\mu^{(\omega)}$ field) is equivalent to quark-interchange. The degree of double counting depends on the form-factor used in the meson exchange part. One way to avoid this double counting is simply to use a form-factor that cuts off the very short distance part of ω exchange, and also provides a smooth matching to the quark-interchange contribution. Such a procedure was used in Ref. [18] for the NN interaction. There,

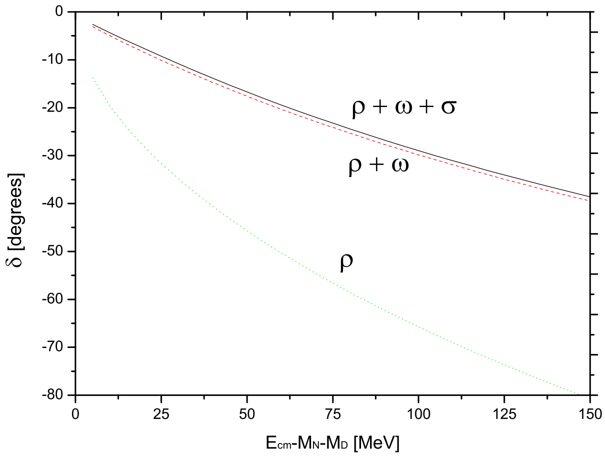

Fig. 5. Meson-exchange contributions to the *s*-wave phase-shift for isospin $I = 0$.

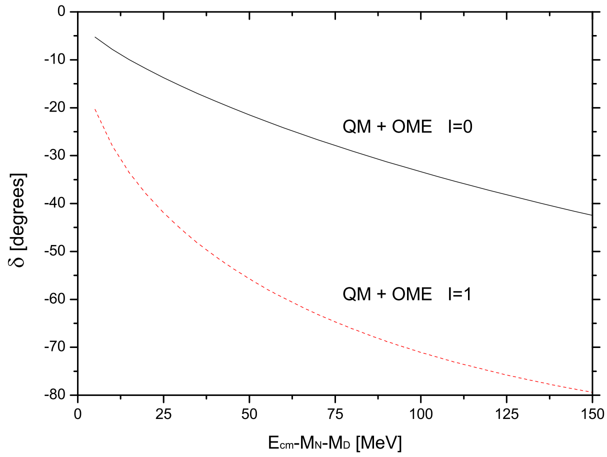

Fig. 7. Quark-interchange (QM) and meson-exchange (OME) contributions to the *s*-wave phase-shift for isospin $I = 0$ and $I = 1$.

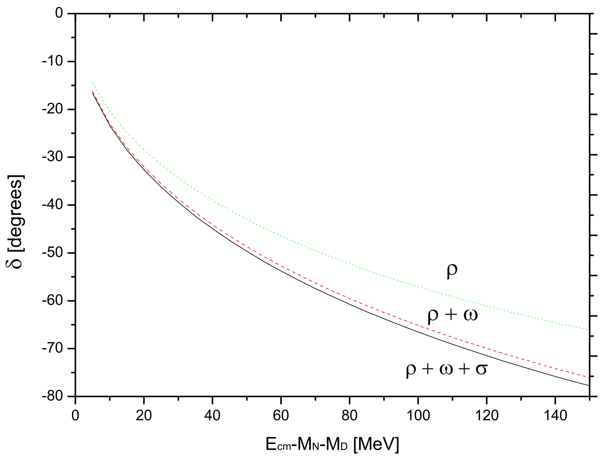

Fig. 6. Meson-exchange contributions to the *s*-wave phase-shift for isospin $I = 1$.

the important point was that although ω exchange is short ranged, its tail plays an important rọle in that reaction. A similar effect is expected for the $\bar{D}N$ reaction, as it was for the K^+N reaction. In the present work we have not attempted to make such a smooth matching, but certainly when cutting off the very short-distance part of the effect of the $\phi_0^{(\omega)}$ field, the results might change. We intend to investigate such effects in a future publication.

6 Conclusions and Perspectives

We have discussed the implementation of a quark model for calculating meson-baryon interactions. The model is based on a Hamiltonian inspired in the QCD Hamiltonian in Coulomb gauge, where gluon fields are eliminated in favor of a confining Coulomb potential and a transverse-gluon potential. The Hamiltonian confines color and realizes DχSB. Color confinement in the model means that only color-singlet states have finite energy.

We use this model to study the phase shifts for elastic process involving the the channels $p\bar{D}^0 \to p\bar{D}^0$, $nD^- \to nD^-$, $pD^- \to pD^-$, $n\bar{D}^0 \to n\bar{D}^0$, $pD^- \to n\bar{D}^0$. We discussed results for the *s*-wave phase-shifts for the coupled states of isospin $I = 0$ and $I = 1$. Our results are considerably more repulsive than the ones using the non-relativistic quark model with one-gluon-exchange. The explanation for the larger repulsion is due to the smaller sizes of hadrons in the present model. Smaller hadron sizes favor larger overlap between the wave-functions of the colliding hadrons and this gives a larger quark-interchange potential. In addition, the microscopic quark interactions used in the present work are stronger than one-gluon-exchange, a feature that also increases the strength of the quark-interchange potential. We have also included meson-exchanges to describe the long-range tail of the interaction. We used SU(4) symmetry to fix the coupling constants. Since we have used the same coupling constants, meson masses, and cutoff masses as in Ref. [2], our results are not different from this reference.

There is plenty of room for improvements in the present approach. Certainly a better determination of hadron sizes should be attempted. One way to improve on this, without sacrificing numerical simplicity, is to use an expansion in terms of several Gaussians for the amplitudes ψ and ϕ, and then diagonalize the resulting Hamiltonian matrix together with a variational determination of the size parameters of the Gaussians. Another improvement is to use the full mass function M_k, instead of the low-momentum expansion of Eq. (26). However, this will have an impact on the numerics, since in this case multidimensional integrals for the determination of the hadron sizes and the effective meson-baryon need to be performed with Monte Carlo integration.

In conclusion, we believe that models like the one presented here are important for improving on our knowledge on low-energy reactions of charmed hadrons with nucleons. There is no experimental information on these reactions and any new development can be of help in guiding future experiments like the ones at the FAIR facility [21]. In addition, a quark model that is based on a Hamiltonian that realizes DχSB should be an improvement on the traditional quark models where there is no link between constituent-quark masses and microscopic interactions. A model that makes such a link is§ much better suited for studies of chiral symmetry restoration in medium.

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A The interaction Lagrangians

In the present work we use the following effective interacting meson-meson and meson-baryon Lagrangian densities[19, 20]

$$\mathcal{L}_{NN\sigma}(x) = g_{NN\sigma}\, \bar{\psi}^{(N)}(x)\, \phi^{(\sigma)}(x)\, \psi^{(N)}(x), \qquad (49)$$

$$\mathcal{L}_{\bar{D}\bar{D}\sigma}(x) = g_{\bar{D}\bar{D}\sigma}\, \varphi^{(\bar{D})}(x)\, \phi^{(\sigma)}(x)\, \varphi^{(D)}(x), \qquad (50)$$

$$\mathcal{L}_{NN\omega}(x) = g_{NN\omega}\Big[\bar{\psi}^{(N)}(x)\,\gamma^{\mu}\,\psi^{(N)}(x)\,\phi^{(\omega)}_{\mu}(x) + \left(\frac{\kappa_\omega}{2M_N}\right)\bar{\psi}^{(N)}(x)\,\sigma^{\mu\nu}\,\psi^{(N)}(x)\left(\partial_\nu\phi^{(\omega)}_{\mu}(x)\right)\Big], \qquad (51)$$

$$\mathcal{L}_{\bar{D}\bar{D}\omega}(x) = ig_{\bar{D}\bar{D}\omega}\Big[\varphi^{(\bar{D})}(x)\,\left(\partial_\nu\,\varphi^{(D)}(x)\right)\,\phi^{(\omega)^{\nu}}(x) - \left(\partial_\nu\varphi^{(D)}(x)\right)\,\varphi^{(\bar{D})}(x)\Big], \qquad (52)$$

$$\mathcal{L}_{NN\rho}(x) = g_{NN\rho}\Big[\bar{\psi}^{(N)}(x)\,\gamma^{\mu}\,\psi^{(N)}(x)\,\tau\cdot\phi^{(\rho)}_{\mu}(x) + \left(\frac{\kappa_\rho}{2M_N}\right)\bar{\psi}^{(N)}(x)\,\sigma^{\mu\nu}\,\psi^{(N)}(x)\,\tau\cdot\partial_\nu\phi^{(\rho)}_{\mu}(x)\Big], \qquad (53)$$

$$\mathcal{L}_{\bar{D}\bar{D}\rho}(x) = ig_{\bar{D}\bar{D}\omega}\Big[\varphi^{(\bar{D})}(x)\,\left(\partial_\nu\varphi^{(D)}(x)\right)\tau - \left(\partial_\nu\varphi^{(\bar{D})}(x)\right)\,\tau\varphi^{(D)}(x)\Big]\cdot\phi^{(\omega)^{\nu}}(x). \qquad (54)$$

In these, $\psi^{(N)}$ denotes the nucleon doublet, $\phi^{(D)}(x)$ the charmed-mesons doublet, $\phi^{(\rho)}_{\mu}(x)$ the iso-triplet of ρ mesons, and τ are the Pauli matrices.

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