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On recent chiral baryon-baryon models*

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I. INTRODUCTION

Chiral $SU(2)_L \times SU(2)_R$ symmetry is presently understood as originating from the small masses of the u and d quark, i.e. as an approximate symmetry of the underlying QCD Lagrangian. Therefore, for any model it is important to what extend it is compatible with (broken) chiral symmetry (CS). In this respect we make some remarks on the Nijmegen soft-core OBE-model [1,2], henceforth called SC-model. In the SC-model we employ the derivative pion-nucleon coupling, i.e. $\gamma_\mu\gamma_5$. This means that in pion-nucleon scattering there is virtually no contribution from the $N\bar{N}$-pair terms. The strong ϵ-exchange contribution to the pion-nucleon scattering length a_0^0 can be canceled completely by the pomeron-exchange contribution. So, with the derivative pion-nucleon coupling and the pomeron contribution the soft-pion theorems [3] can be satisfied [4]. The $\gamma_\mu\gamma_5$-coupling is characteristic for the non-linear realization of chiral-symmetry. In that case there is no need for (fictiteous) σ. Henceforth we will therefore only discuss models based on this non-linear realization, like the SC-model. (Notice, the broad $\epsilon(760)$-meson as employed in the Nijmegen OBE-models, is not to be confused with the σ.)

Going beyond OBE, including two-pion exchange etc., one has to include also meson-pair vertices, which occur naturally in the non-linear chiral invariant Lagrangians. In this paper we will discuss two recent models on the baryon-baryon interactions, where strict constraints from CS are imposed.

a. Chiral-Perturbation Model. In chiral-perturbation-theory (CPTh) the amplitudes are expanded in power series of f_π^{-1}. For this the non-linear realization of chiral-symmetry is very convenient, since it leads automatically to such an expansion. This in contrast to the linear realization. Recently, Ordóñez, Ray, and van Kolck [5] presented a nucleon-nucleon (NN) potential based on an effective chiral Lagrangian of pions, nucleons, and Δ isobars. Using 26 free parameters, the agreement with the experimental scattering data was found to be satisfactory up to lab energies of about 100 MeV. An extension to higher energies and a further improvement in the description of the data would require an expansion to higher orders in chiral perturbation theory, making the model much more complicated and introducing many new parameters. This makes the CPTh-approach not very attractive for the investigation of baryon-baryon interactions and nuclear and hyper-nuclear systems. The extension to $SU(3) \times SU(3)$ has not been attempted so far. Also, in this case the convergence of the CPTh-expansion is expected to be worse than in the NN case.

b. Chiral ESC-Model. In [7] we give a chiral-invariant version of the ESC-model [8–10]. In contrast to the procedure in the CPTh-models, like [5], here all mesons other than the pion [6] are not 'integrated out'. Instead, the succesful approach of the OBE-models is used and all lowest-lying mesons with masses lower than 1 GeV are included. The BB potential model is then obtained by evaluating the standard one-boson-exchange contributions involving these mesons, but now including the contributions of the box and crossed-box two-meson diagrams and of the pair-meson diagrams where at least one of the baryon lines contains a pair-meson vertex [10]. We stress that we do not use CPTh, but we employ CS to generate the $\mathcal{L}_I$ and to implement constraints for the associated coupling constants. The pair-meson interactions of the ESC-model, involving pions, arise as a direct consequence of CS. In Refs. [10] we already showed that the inclusion of the two-meson (box, crossed-box, and pair) contributions provides a substantial improvement in the description of the NN-data as compared to a potential containing only the standard one-boson exchanges. This important result motivates to go beyond the NN model and to investigate whether a similar approach will also be fruitful in the construction of an extended hyperon-nucleon (YN) potential. An important motivation for the development of an extended YN model is provided by the study of hypernuclei using one-boson-exchange models, see e.g. [11,12]. For that purpose, we discussed in [7] the non-linear realization of CS to $SU(3)_L \times SU(3)_R$ to obtain a CS meson-baryon interaction Lagrangian. This $SU(3)$-invariant $\mathcal{L}_I$ describes the coupling of the pseudo-scalar, scalar, vector and axial-vector mesons with the baryon-octet containing the N, Λ, Σ, and Ξ.

The content of this paper is as follows. In section II a short outline of the $SU(3) \times SU(3)$ construction is given. In section III a brief description is given of the estimates of the single-meson couplings. In section IV the procedure and results for the meson-pair vertices are reviewed. In section V we discuss the results for NN and the prospects for the applications to YN and YY.

II. SU(3) CHIRAL SYMMETRY

In [7] we have given an $SU(3)$ manifest chiral invariant formulation of an ESC-model, which is a non-linear extension of the linear $SU(2) \times SU(2)$ σ-model of Ko and Rudaz [13]. Here, we give an outline of this model. For a general review on CS we refer to [14,15]. The baryon-octet matrix Ψ left- and right-handed components, $\Psi_{L,R} = \frac{1}{2}(1 \mp \gamma_5)\psi$, transform as

$$\Psi_L \to L\Psi_L L^\dagger, \qquad \Psi_R \to R\Psi_R R^\dagger. \tag{1}$$

In the linear representation one has an $SU(3)$ singlet scalar σ, and a pseudo-scalar octet π_a $(a = 1, \ldots, 8)$. One finds for

$$\Sigma = \sigma + i\lambda_a \pi_a, \quad (a = 1, \ldots, 8), \tag{2}$$

where λ_a are the Gell-Mann matrices, that Σ transforms under global $SU(3)_L \times SU(3)_R$ as $\Sigma \to L\Sigma R^\dagger$, where L and R are elements of SU(3).

The transformation to the non-linear formulation is effected by introducing

$$\left.\begin{array}{l} B_R = u\Psi_R u^\dagger \\ B_L = u^\dagger \Psi_L u \end{array}\right\}, \quad u(\xi_a) = \exp[i\lambda_a \xi_a] \equiv \exp\left[\frac{i\lambda_a \pi'_a}{2f_0}\right], \tag{3}$$

where $f_0 = \langle\sigma\rangle$. (In the absence of vector and axial-vector gauge fields $f_0 = f_\pi$.) Note that $u(\xi_a)$ are elements of the coset space $\mathrm{SU}(3)_L \times \mathrm{SU}(3)_R/\mathrm{SU}(3)_V$. The π'_a-fields are identified with the octet of physical pseudoscalar fields:

$$\frac{1}{\sqrt{2}}\lambda_a \pi'_a = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}} & K^0 \\ K^- & \overline{K^0} & -\frac{2\eta_8}{\sqrt{6}} \end{pmatrix}, \tag{4}$$

The octet of baryon fields reads

$$B = \begin{pmatrix} \frac{\Sigma^0}{\sqrt{2}} + \frac{\Lambda}{\sqrt{6}} & \Sigma^+ & p \\ \Sigma^- & -\frac{\Sigma^0}{\sqrt{2}} + \frac{\Lambda}{\sqrt{6}} & n \\ \Xi^- & -\Xi^0 & -\frac{2\Lambda}{\sqrt{6}} \end{pmatrix}, \tag{5}$$

which transforms as $B \to HBH^\dagger$. with

$$H = \sqrt{Lu^2R^\dagger}\; Ru^\dagger \;=\; \sqrt{Ru^{\dagger\,2}L^\dagger}\; Lu \;\in SU(3), \tag{6}$$

and $u \to LuH^\dagger = HuR^\dagger$. The covariant derivative of the baryon octet reads

$$D_\mu B = \partial_\mu B + i[\Gamma_\mu, B]. \tag{7}$$

where the connection is $\Gamma_\mu = -\frac{i}{2}\left(u^\dagger\partial_\mu u + u\partial_\mu u^\dagger\right)$.

The field combinations $u\Sigma^\dagger u$ and $u^\dagger\Sigma u^\dagger$ both transform according to the baryon octet matrix B. So, defining $\chi_\pm = \frac{1}{2}\left(u^\dagger\Sigma u^\dagger \pm u\Sigma^\dagger u\right)$, the simplest chiral invariant interaction Lagrangian is given by

$$\mathcal{L}_I = -g_{s,1}\,\mathrm{Tr}(\overline{B}\chi_+ B) - g_{s,2}\,\mathrm{Tr}(\overline{B}B\chi_+), \tag{8}$$

where $g_{s,1}$ and $g_{s,2}$ are arbitrary constants.

However, this Lagrangian gives for all baryons in the octet the same mass, $M = (g_{s,1} + g_{s,2})f_0$. In order to generate the thrical baryon masses, an octet of scalar fields, $\lambda_a\sigma_a$, is added with a non-vanishing vacuum expectation value for the isoscalar octet member. To incorporate this we write

$$\Sigma = F + \lambda_0 s_0 + \lambda_a s_a + i\lambda_a\pi_a, \qquad (a = 1, \ldots, 8), \tag{9}$$

where π_a are the *original* pseudoscalar fields and F the vacuum expectation value of the scalar-nonet fields [7]. It can be shown that it is still possible to find a transformation $u(\xi)$, which transforms away the octet of original pseudoscalar fields, while leaving the matrix F invariant, and where the ξ_a fields can be identified with an octet of new pseudoscalar fields. As before, the new $\chi_\pm$ are defined in terms of the new u and Σ. However, $\chi_\pm$ contain the *original* pseudoscalar fields in a complicated way. But since χ_+ behaves like a set of scalar fields, one can simply *define* these new fields to be the physical scalar fields and any reference to the original scalar fields can be dropped. Omitting primes, the nonet of the new scalar fields is given by

$$\chi_+ = F + \lambda_0 s_0 + \lambda_a s_a, \tag{10}$$

where s_0 now denotes the new scalar singlet and the octet matrix is given by $\frac{1}{\sqrt{2}}\lambda_a s_a$, which is similar to 4 but with $\boldsymbol{\pi} \to \mathbf{a}_0, \eta_8 \to s_8$, and $K \to \kappa$. The octet isoscalar s_8 and singlet s_0 are mixed to give the physical $f_0(980)$ and $\varepsilon(760)$. Also, $-i\chi_-$ can be identified as a new isosinglet pseudoscalar field, not present before. Because it transforms as $H(-i\chi_-)H^\dagger$, it can be formally added to the octet pseudoscalar matrix, which completes the nonet. It must be remembered, however, that group transformations are only valid for the (traceless) octet matrices.

The next step is to extend the global $SU(3) \times SU(3)$ to a local chiral symmetry. The required gauge fields are given by two octets of combinations of vector and axial-vector fields. The vector octet is $\frac{1}{\sqrt{2}}\lambda_a \rho_a$, again similar to 4 but now with $\boldsymbol{\pi} \to \boldsymbol{\rho}$, $\eta_8 \to \omega_8$, and $K \to K^*$, and analogous for the axial-vector octet. For the details of this extension we refer to the forthcoming paper [7]. Also, there one can find a detailed discussion of the description of the meson masses in the context of this chiral model.

The main interest for this paper is the discussion of the baryon-baryon interaction in the chiral framework. The main result of the description of the physical meson fields in the chiral context is that we have arrived at various 3×3-matrices of the general form $\Phi = \frac{1}{\sqrt{2}}\lambda_c \phi_c$. These matrices contain scalar, pseudo-scalar, vector, and axial-vector fields, which, except for the vector fields, all transform similarly to the baryon fields. This fact leads to the following chiral-invariant combinations:

$$\begin{aligned}
\left[\overline{B}B\Phi\right]_F &= \mathrm{Tr}(\overline{B}\Phi B) - \mathrm{Tr}(\overline{B}B\Phi), \\
\left[\overline{B}B\Phi\right]_D &= \mathrm{Tr}(\overline{B}\Phi B) + \mathrm{Tr}(\overline{B}B\Phi) - \tfrac{2}{3}\,\mathrm{Tr}(\overline{B}B)\mathrm{Tr}(\Phi), \\
\left[\overline{B}B\Phi\right]_S &= \mathrm{Tr}(\overline{B}B)\mathrm{Tr}(\Phi).
\end{aligned} \tag{11}$$

A general interaction Lagrangian, which satisfies chiral-symmetry is the $SU_f(3)$-invariant

$$\mathcal{L}_I = -g^{\mathrm{oct}}\sqrt{2}\left\{\alpha\left[\overline{B}B\Phi\right]_F + (1-\alpha)\left[\overline{B}B\Phi\right]_D\right\} - g^{\sin}\sqrt{\tfrac{1}{3}}\left[\overline{B}B\Phi\right]_S, \tag{12}$$

where α is the $F/(F+D)$ ratio. The meson field matrices are given by

$$\Phi_{\mathrm{sc}} = \frac{1}{\sqrt{2}}\left[F + \lambda_c s_c\right], \tag{13}$$

$$\begin{aligned}
\Phi_{\mathrm{vc}} = \frac{-i}{\sqrt{2}g_V}\gamma_\mu \Big[& u^\dagger\left(\partial^\mu + \tfrac{i}{2}g_V\lambda_c(\rho^\mu + A^\mu - hD^\mu\pi)_c\right)u \\
& + u\left(\partial^\mu + \tfrac{i}{2}g_V\lambda_c(\rho^\mu - A^\mu + hD^\mu\pi)_c\right)u^\dagger\Big],
\end{aligned} \tag{14}$$

$$\begin{aligned}
\Phi_{\mathrm{ax}} = \frac{-i}{\sqrt{2}g_V}\gamma_5\gamma_\mu \Big[& u^\dagger\left(\partial^\mu + \tfrac{i}{2}g_V\lambda_c(\rho^\mu + A^\mu - hD^\mu\pi)_c\right)u \\
& - u\left(\partial^\mu + \tfrac{i}{2}g_V\lambda_c(\rho^\mu - A^\mu + hD^\mu\pi)_c\right)u^\dagger\Big],
\end{aligned} \tag{15}$$

where $D^\mu(\lambda\pi) = \partial^\mu(\lambda\pi) - \frac{i}{2}g_V[(\lambda\pi),(\lambda\rho_\mu)]$, and h chosen such that the mixing between the axial-vector and pseudoscalar fields in the meson sector vanishes. Note that the pseudovector coupling of the pseudoscalar fields is already included in the axial-vector field matrix Φ_{ax}. $\Phi_{\mathrm{ax}} \equiv \gamma_5\gamma^\mu a'_\mu$.

In addition to the electric coupling $\Phi_{vc} \equiv \gamma^\mu \rho'_\mu$, it is also possible [16] to include a chiral-invariant magnetic coupling $\sigma^{\mu\nu}\rho'_{\mu\nu}$, where $\rho'_{\mu\nu}$ is the field strength tensor for the ρ'_μ field combination. This is due to the fact that ρ'_μ transforms similar to Γ_μ, and so we can define a chiral-invariant field strength tensor $\rho'_{\mu\nu}$. Furthermore, the transformation properties of Γ_μ also imposes the constraint that the chiral-variant D-type coupling $[\overline{B}B\Phi_{\rm vc}]_D$ should vanish, i.e., the electric $\alpha_V^e = 1$. Hence, the assumption that the ρ meson couples universally [17] to the isospin current in this model is a direct consequence of chiral SU(3) symmetry. The magnetic α_V^m is still a free parameter.

Notice that the $SU(3)_L \times SU(3)_R$ invariant interaction $\mathcal{L}_I$ contains through the structure of the Φ_{sc}, Φ_{vc}, and Φ_{ax}, not only the single meson-baryon vertices, leading to OBE, TME, etc., but also multi-meson couplings to baryons. The couplings for all these multi-meson interactions can all be expressed in terms of the single-meson interaction coupling constants. Also, the couplings of the axial-vector mesons are all directly related to those of the pseudoscalar mesons.

The couplings of the single-meson vertices have all been discussed in detail in [7]. In the chiral ESC-model [7] we include, as explained in the Introduction, one-boson-exchange, two-meson exchange, and meson-pair exchanges that are generated by the chiral-invariant interactions described above. In confronting this chiral ESC-model with the NN-data, this fixing of all couplings is too restrictive. Below we will describe how more flexibility can be achieved, still within the framework of chiral symmetry.

Note that the 3×3 matrices Φ are not simple representations of the scalar, vector, or axial-vector meson fields, but they contain the pseudoscalar fields in a nonlinear way as well. This means that the chiral Lagrangian contains all kinds of multiple-meson (pair, triple, etc.) interactions not envisaged before. The coupling constants for these multiple-meson interactions can all be expressed in terms of the single-meson interaction coupling constants. This will be the subject of Sec. III.

III. MESON-PAIR INTERACTIONS

Originally, with the ESC-model [8,9] the meson-pair interaction couplings were treated as free parameters in order to investigate the possibilities to improve the NN-fit. The implementation of the restrictions from e.g. CS were envisaged for a later stage of an approach where besides OBE and TME also meson-pair exchange is exploited. In [10] it was shown that the inclusion of chiral constraints, as implied by the linear σ-model, could be imposed indeed. This, without impeding a substantial improvement in the description of the NN-data as compared to a potential containing only OBE. In this section we discuss the double-meson vertices as are characteristic for the chiral model as described in [7].

a. Vector double-meson vertices. Expanding to second order in the meson fields

$$\Phi_{\rm vc} = \frac{1}{\sqrt{2}}\,\gamma^\mu(\lambda\rho_\mu) - \frac{i(1-2g_V f_1 h)}{4\sqrt{2}g_V f_1^2}\,\gamma^\mu\left[(\lambda\pi),\partial_\mu(\lambda\pi)\right] - \frac{i}{2\sqrt{2}f_1}\,\gamma^\mu\left[(\lambda\pi),(\lambda A_\mu)\right] + \dots \quad (16)$$

The pair interaction Lagrangian is obtained by the replacement

$$\gamma^\mu \boldsymbol{\rho}_\mu \longrightarrow \gamma^\mu(\boldsymbol{\pi}\times\partial_\mu\boldsymbol{\pi}),\ \text{etc. }, \quad (17)$$

in the single-meson Lagrangian, which gives

$$
\begin{aligned}
m_\pi^2 \mathcal{L}_{(\pi\pi)} = & -g_{NN(\pi\pi)}(\overline{N}\gamma^\mu \boldsymbol{\tau} N)\cdot(\boldsymbol{\pi}\times\partial_\mu\boldsymbol{\pi}) - g_{\Xi\Xi(\pi\pi)}(\overline{\Xi}\gamma^\mu\boldsymbol{\tau}\Xi)\cdot(\boldsymbol{\pi}\times\partial_\mu\boldsymbol{\pi}) \\
& -g_{\Lambda\Lambda(\pi\pi)}(\overline{\Lambda}\gamma^\mu\boldsymbol{\Sigma}+\overline{\boldsymbol{\Sigma}}\gamma^\mu\Lambda)\cdot(\boldsymbol{\pi}\times\partial_\mu\boldsymbol{\pi}) \\
& +ig_{\Sigma\Sigma(\pi\pi)}(\overline{\boldsymbol{\Sigma}}\times\gamma^\mu\boldsymbol{\Sigma})\cdot(\boldsymbol{\pi}\times\partial_\mu\boldsymbol{\pi}).
\end{aligned} \tag{18}
$$

Here we introduced the square of the charged-pion mass to make the coupling constants dimensionless. Substituting the appropriate renormalization factors and $\frac{1}{2}g_V = g_{NN\rho}$, This leads to the coupling constants [7]

$$
g_{B'B(\pi\pi)} = \frac{m_\pi^2}{4f_1^2}\,\frac{2Z_\pi - 1}{Z_\pi}\,\frac{g_{B'B\rho}}{g_{NN\rho}}, \tag{19}
$$

for $B'B = NN$, $\Xi\Xi$, $\Lambda\Sigma$, and $\Sigma\Sigma$. Here, $g_{NN\rho} = \frac{1}{2}g_V$, and Z_π is given by $g_{NNa_1} = (m_{a_1}/m_\pi)\, f_{NN\pi}\sqrt{(1-Z_\pi)/Z_\pi}$. Note that by choosing [18,16,19] $Z_\pi = \frac{1}{2}$, i.e., making the assumption that $m_{a_1} = \sqrt{2}m_\rho$, all the $(\pi\pi)$ pair interactions are absent. However, experimentally $m_{a_1} \neq \sqrt{2}m_\rho$, and so here the $(\pi\pi)$ pair interactions are still present in the interaction Lagrangian. For further explicit results for the (πK), $(\eta_8 K)$, and $(K\bar{K})$ we refer again to [7]. Here, also the pseudo-scalar and axial-vector double meson couplings are given.

b. Axial-vector double-meson vertices. Expansion of Φ_{ax} to second order in the meson fields gives

$$
\Phi_{\rm ax} = \frac{1}{\sqrt{2}}\gamma^5\gamma^\mu(\lambda A_\mu) + \frac{1-g_V f_1 h}{\sqrt{2}g_V f_1}\gamma^5\gamma^\mu\partial_\mu(\lambda\pi) - \frac{i(1-g_V f_1 h)}{2\sqrt{2}f_1}\gamma^5\gamma^\mu\left[(\lambda\pi),(\lambda\rho_\mu)\right] + \ldots \tag{20}
$$

In complete analogy with the previous, making the substitutions

$$
\gamma^5\gamma^\mu\mathbf{A}_\mu \longrightarrow \gamma^5\gamma^\mu(\boldsymbol{\pi}\times\boldsymbol{\rho}_\mu),\ \text{etc.}
$$

in the single-meson $\mathcal{L}_I$ gives the meson-pair interactions. The coupling constants are most easily expressed in terms of $g_{NN\rho}$ and the pseudovector coupling constants:

$$
g_{B'B(\pi\rho)} = 2g_{NN\rho}f_{B'B\pi}, \tag{21}
$$

and similar expressions for πK^*, $\eta_8 K^*$, and KK^*, see [7].

c. Quadratic Φ_{sc} and Φ_{ax} extensions. So far, all couplings are fixed by empirical constraints, in the case of the single-meson vertices, and by chiral symmetry, in the case of the pair-vertices. As mentioned before, in order to make a succesfull fit to the NN-data, some flexibility is needed, however. Note, that because of the constraints from the baryon mass generation, the scalar couplings are at this moment completely fixed, see [7]. In view of the transformation properties of the $\Phi_{\rm sc}$ and the $\Phi_{\rm ax}$ fields, we can add more interaction terms quadratically etc. in the meson fields Φ. In order to create some freedom in the scalar couplings we extend $\mathcal{L}_I(\Phi_{\rm sc})$ by the substitution

$$
\Phi_{\rm sc} \Rightarrow \Phi_{\rm sc} + \sqrt{2}\left(\frac{g_{ss}}{m_\pi}\right)(\Phi_{\rm sc})^2 \tag{22}
$$

This, since both $\Phi_{\rm sc} \rightarrow H\Phi_{\rm sc}H^\dagger$ and $\Phi_{\rm sc}^2 \rightarrow H\Phi_{\rm sc}^2 H^\dagger$ under a chiral transformation. Hence, writing $X = \lambda_c s_c$ we have the substitution

$$g_{\rm sc}^{\rm oct}\Phi_{\rm sc}^{(8)} \longrightarrow \frac{1}{\sqrt{2}}\left\{g_{\rm sc}^{\rm oct}\left(F+X\right)+\frac{g_{ss}^{(8)}}{m_\pi}\left(F^2+(FX+XF)+X^2\right)\right\}, \tag{23}$$

and a similar expression for the singlet part. In [7] it is shown that this extension leads to the desired freedom in the scalar meson couplings, and at the same time the proper description of the empirical baryon masses can be maintained.

Finally, in [7] the model is extended by including interactions of the type $\Phi_{\rm ax}^2$, or equivalently $u_\mu u_\nu$, where $u_\mu = -\frac{i}{2}\left(u^\dagger\partial_\mu - u\partial_\mu u^\dagger\right)$. Distinguishing the symmetric and the anti-symmetric part

$$\begin{aligned}\phi_s &\sim -\tfrac{1}{2}g^{\mu\nu}\left[\partial_\mu(\lambda\pi)\partial_\nu(\lambda\pi)+\partial_\nu(\lambda\pi)\partial_\mu(\lambda\pi)\right],\\ \phi_a &\sim +\tfrac{i}{2}\sigma^{\mu\nu}\left[\partial_\mu(\lambda\pi)\partial_\nu(\lambda\pi)-\partial_\nu(\lambda\pi)\partial_\mu(\lambda\pi)\right].\end{aligned}$$

extra two-pseudo-scalar contributions to the meson-pair vertices are obtained, see [7]. From these only the double pseudo-scalar terms are included in the actual fit to the NN-data. The free papRameters for these new terms are $g_{sym}^{(8)}, g_{sym}^{(1)}$, and α_{sym} for the symmetric case, and $g_{asym}^{(8)}, g_{asym}^{(1)}$, and α_{asym} for the antisymmetric case.

IV. APPLICATION TO *NN* AND PROSPECTS

In [7] we have given a first application of the manifest CS ESC-model. We imposed all the constraints from CS and showed that the resulting NN model gives a very reasonable fit to the NN-data. In this model the treatment of the meson-exchanges is the same as in the Nijmegen soft-core models [1,2].

The 12 free parameters of the model (Λ_S, Λ_P, Λ_V, g_{A_2}, g_P, $g_{ss}^{(8)}$, $g_{ss}^{(1)}$, $g_{sym}^{(8)}$, $g_{sym}^{(1)}$, α_{sym}, $g_{asym}^{(8)}$, and α_{asym}) were determined in a fit to the Nijmegen representation [20] of the χ^2 hypersurface of the *NN* scattering data below $T_{\rm lab} = 350$ MeV, updated with the inclusion of new data which have been published since then. The reached χ^2-per-datum for each of the ten energy bins is shown in Table I, in comparison with the (updated) Nijmegen partial-wave analysis. The model gives a good description of the data below 300 MeV (χ^2-per-datum of 1.5), but rapidly worsens at higher energies. This is probably due to the nonadiabatic expansion in the two-meson contributions [10] which, strictly speaking, is only valid below the pion production threshold ($T_{\rm lab} \approx 280$ MeV). However, the χ^2-per-datum of 1.9 for the 0–350 MeV energy interval is still very acceptable in comparison to other potential models that have appeared in the literature [21]. Furthermore, it should be realized that in this model *all* coupling constants satisfy constraints as imposed by chiral symmetry, or empirical constraints, see [7]. This in contrast to any other model that has appeared in the literature.

The application to the YN and YY channels is in principle straightforward, and can be done with only a few free parameters, like the mixing angle in the scalar nonet. This will be useful for the study of hypernuclei, in order to have a variety of theoretical models at disposal. This is important to investigate the sensitivity of the hypernuclear systems to particular types of interaction, like the spin-spin, tensor, and spin-orbit interaction.

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TABLES

TABLE I. χ^2 and χ^2 per datum ($\chi^2_{\text{p.d.p.}}$) at the 10 energy bins for the updated partial-wave analysis (PWA) and the constrained NN potential. N_{data} lists the number of data within each energy bin. The bottom line gives the results for the total 0–350 MeV interval.

		PWA		potential	
Bin(MeV)	N_{data}	χ^2	$\chi^2_{\text{p.d.p.}}$	χ^2	$\chi^2_{\text{p.d.p.}}$
0.0–0.5	145	144.45	0.996	152.40	1.05
0.5–2	68	42.97	0.632	54.52	0.80
2–8	110	106.28	0.966	186.94	1.70
8–17	296	276.31	0.933	350.87	1.19
17–35	359	279.54	0.829	376.91	1.05
35–75	585	567.18	0.970	963.92	1.65
75–125	399	409.58	1.027	483.10	1.21
125–183	760	820.69	1.080	1473.04	1.94
183–290	1047	1035.48	0.989	1676.26	1.60
290–350	992	997.02	1.005	3241.95	3.27
0–350	4761	4697.50	0.987	8959.91	1.88