Abstract

In the supersymmetric models, the dominant sources of the hadronic flavor-diagonal CP violation at low energy are the theta term and the chromoelectric dipole moments of quarks. Using QCD sum rules, we estimate the preferred range and the best values for the CP-odd meson-nucleon coupling constants induced by these operators. When the theta term is removed by the axion mechanism, the size of the most important isospin triplet pion-nucleon coupling is estimated to be $g^{(1)}_{\pi NN} = 2 \times 10^{-12}(\bar d_u - \bar d_d)$, where chromoelectric dipole moments are given in units of $10^{-26}$cm.
The search for CP violation in flavor-conserving processes is of paramount scientific importance. The suppression of CP-violating effects induced by the complex phase of the Kobayashi-Maskawa matrix allows the use of electric dipole moments (EDMs) of neutrons or heavy atoms as well as T-odd asymmetries in the decays and scattering of baryons as powerful tools for probing new physics beyond the Standard Model.

The wide separation between the energy scale of “new physics” (superpartners, technicolor, etc.) and the characteristic momenta of particles in non-accelerator experiments permits consideration of only the first few terms in the effective CP-odd Lagrangian. In the minimal supersymmetric models only the theta term, three-gluon operator, and EDMs and color EDMs of light quarks are important:

$$\delta \mathcal{L} = - \sum_{q=u,d,s} m_q \bar{q}(1 + i \theta_q \gamma_5)q + \frac{\alpha_s}{8\pi} G \tilde{G} + w GG \tilde{G} - \frac{i}{2} \sum_{q=u,d,s} d_q \bar{q} F \gamma_5 q - \frac{i}{2} \sum_{q=u,d,s} \bar{d}_q q g_s G \gamma_5 q,$$

where $G \tilde{G} \equiv G_{\mu\nu}^{a} \tilde{G}_{\mu\nu}^{a}$, $G \sigma \equiv t^a G_{\mu\nu}^{a} \sigma_{\mu\nu}$ and $GG \tilde{G} \equiv f^{abc} G_{\mu\nu}^{a} G_{\nu\alpha}^{b} \tilde{G}_{\alpha\mu}^{c}$. The coefficients in (1) are generated by the CP violation in the SUSY breaking sector and evolved down to 1 GeV, which is the border line of viability of the perturbative quark-gluon description.

In this Letter we present a systematic study of the transition from this Lagrangian to the effective T-odd meson-nucleon interactions which determine the magnitude of the CP-violating nuclear moments and the T-odd asymmetries in nucleon scattering. T-odd nuclear forces are the main source for the EDMs of heavy diamagnetic atoms (see e.g. [1]). The quality of constraints imposed on supersymmetric models from a recently improved measurement of the EDM of the xenon and mercury atoms [2, 3] (as well as of future experimental efforts with the EDMs of diamagnetic atoms and the T-odd nucleon scattering [4]) depends crucially on the treatment of QCD and nuclear effects, i.e. on the extraction of limits on $\tilde{d}_i$ from the experimental bound on the atomic EDM. The implications of this powerful constraint for the CP violation in the supersymmetric models have been emphasized in Refs. [5, 6], and numerically exploited in [7]. The purpose of this work is to give “state-of-the-art” estimates for various T-odd nucleon-meson coupling constants, i.e to find their best values in terms of the coefficients in eq.
This problem is reminiscent of Ref. [8] which estimates various P-odd meson-nucleon couplings in terms of the Fermi constant.

T-odd nuclear forces inside the nucleus can be approximated by a meson exchange with one of the meson-nucleon couplings being T-violating [9]. It is natural to expect that pion exchange dominates in the T-odd channel. The coupling of nucleons with pions can be conveniently parametrized [10, 11] as

$$\mathcal{L}^{CP} = \bar{g}^{(0)}_{\pi NN} \bar{N} \tau^a N \pi^a + \bar{g}^{(1)}_{\pi NN} \bar{N} N \pi^0$$

These couplings are generated by the theta term and by the color EDMs of quarks, $\bar{g} \equiv \bar{g}(\tilde{\theta}, \tilde{d}_i)$, where $\tilde{\theta} = \theta_G + \sum \theta_q$. Couplings that change the isospin by two units can be generated only at the expense of an additional $m_u - m_d$ suppression and are ignored in the present analysis. $\bar{g}^{(0)}_{\pi NN}(\tilde{\theta})$ is rather well known [12], as it can be deduced from the size of the $\langle N | \bar{u}u - \bar{d}d | N \rangle$ matrix element. For most of the models of CP violation including minimal SUSY models, $\bar{g}$ has to be removed by the Peccei-Quinn (PQ) symmetry leaving quark color EDMs as the dominant source for CP-odd nuclear forces. The contribution of the three-gluon operator $GG\bar{G}$ to $\bar{g}$ is additionally suppressed by $m_q$ and can be neglected.

The first step in the calculation of $\bar{g}^{(0)}_{\pi NN}(\tilde{\theta}, \tilde{d}_u + \tilde{d}_d)$ and $\bar{g}^{(1)}_{\pi NN}(\tilde{d}_u - \tilde{d}_d)$ is the reduction of the pion field by means of PCAC [13], Fig. 1a. The smallness of the $t$-channel pion momentum compared to the characteristic hadronic scale justifies this procedure,

$$\langle N | \pi^a | O_{CP} | N' \rangle = \frac{i}{f_\pi} \langle N | [O_{CP}, J^a_5] | N' \rangle.$$  (3)

The commutator of the zero component of the axial current with CP-violating operators $O_{CP} = \bar{q}g_s G\sigma G q$ can be easily computed, leading to the matrix elements of the $\bar{q}g_s G\sigma q$ operators over the nucleon state [13]. However, eq. (3) is an incomplete result. A second class of contributions was pointed out in Ref. [9] and in Refs. [14, 15] in the context of the neutron EDM problem. They consist of the pion pole diagrams, Fig. 1b, which contribute at the same order of chiral perturbation theory. Indeed, the quantum numbers of the $\bar{q}g_s G\sigma q$ operators allow them to produce zero-momentum $\pi^0$s from the vacuum. The pion-nucleon scattering amplitude at vanishing pion momentum is proportional to the first power of the quark mass, whereas the pion propagator contains $1/m_q$, so that Fig. 1b and 1a both contribute at the same $O(1/m_q)$ order.
Using the low energy theorems that relate the pion-nucleon scattering amplitude with the matrix elements of $m_q\bar{q}q$ over the nucleon state, we arrive at the following intermediate result for the $NN\pi^0$ vertex:

\[
\frac{1}{2f_\pi}\langle N|\tilde{d}_u(\bar{u}g_sG\sigma u - m_0^2\bar{u}u) - \tilde{d}_d(\bar{d}g_sG\sigma d - m_0^2\bar{d}d)|N\rangle
\]

\[
+ \frac{m_*}{2f_\pi} \left[ 2\bar{\theta} + m_0^2 \left( \frac{\tilde{d}_u}{m_u} + \frac{\tilde{d}_d}{m_d} + \frac{\tilde{d}_s}{m_s} \right) \right] \langle N|\bar{u}u - \bar{d}d|N\rangle.
\]

In this expression, $m_* = m_u m_d/(m_u + m_d)$ and $m_0^2 = \langle 0|\bar{q}g_sG\sigma q|0\rangle/\langle\bar{q}q\rangle = -(0.8 \pm 0.1)\text{GeV}^2$ [14] parametrizes the strength of the quark-gluon dim=5 vacuum condensate. In our case, this originates from the $\langle \pi_0|\bar{q}g_sG\sigma\gamma_5q|0\rangle$ matrix element and the minus sign is included into the definition of $m_0^2$ for convenience. An alternative way of obtaining the amplitude (4) is to chirally rotate quark masses to the basis where pions cannot be produced from the vacuum, $\langle\pi^0|2\sum q_m\theta_q\gamma_5q + \sum \bar{d}_q\bar{g}_sG\sigma\gamma_5q|0\rangle = 0$, while keeping the theta term fixed, $\sum \theta_q = \text{const}$. This eliminates diagrams 1b, but creates an additional contribution to 1a, leading to the same result (4).

When the PQ mechanism is activated, removing the theta term, the minimum of the axion vacuum is shifted from $\bar{\theta} = 0$ by the color EDM operators [14]. It turns out that the true minimum is such that the square bracket in eq. (4) is zero so that only the first line survives. This leads to a relatively simple expression for the couplings

\[
\bar{g}^{(0)}_{\pi NN} = \frac{\tilde{d}_u + \tilde{d}_d}{2f_\pi} \langle p|H_u - H_d|p\rangle
\]

\[
\bar{g}^{(1)}_{\pi NN} = \frac{\tilde{d}_u - \tilde{d}_d}{2f_\pi} \langle N|H_u + H_d|N\rangle
\]
in terms of matrix elements of the $H_u$ and $H_d$ operators:

$$H_u \equiv \overline{u}g_s G\sigma u - m_0^2\overline{u}u; \quad H_d \equiv \overline{d}g_s G\sigma d - m_0^2\overline{d}d$$  \hspace{1cm} (6)

Previously, using a combination of QCD sum rules and scaling arguments, Ref. \cite{13} estimated that $\langle N|\overline{q}g_s G\sigma q|N\rangle \sim \frac{2}{3}m_0^2\langle N|\overline{q}q|N\rangle$. Another analysis \cite{18} finds similar estimate $\langle N|\overline{q}g_s G\sigma q|N\rangle \sim m_0^2\langle N|\overline{q}q|N\rangle$. Obviously, these estimates are not sufficient to derive a reliable answer for $\overline{g}_Z^{(0)}$ and $\overline{g}_Z^{(1)}$ because of the additional $-m_0^2\langle N|\overline{q}q|N\rangle$ contribution coming from diagrams 1b. The danger of mutual cancelation between the two contributions was realized in Ref. \cite{6} where the need for a dedicated analysis of $\langle N|H_{u(d)}|N\rangle$ was emphasized. In the rest of this paper we derive the QCD sum rules \cite{13} for the matrix elements of the $H_{u(d)}$ operators. The advantage of this approach is that the operator product expansion (OPE) will contain similar vacuum condensates for both sources, $\overline{q}g_s G\sigma q$ and $-m_0^2\overline{q}q$, which allows us to trace possible cancelations. Following Refs. \cite{15, 20}, we introduce the generalized nucleon interpolating current,

$$\eta_p = (j_1 + \beta j_2),$$  \hspace{1cm} (7)

which combines the two Lorentz structures, $j_1 = 2\epsilon_{abc}(u^c_a C\gamma_5 d_b)u_c$ and $j_2 = 2\epsilon_{abc}(u^c_a C d_b)\gamma_5 u_c$.

We compute the OPE for the correlator of this current in the presence of
\[ \tilde{d}_{u(d)} H_{u(d)} \text{ external sources} \]

\[ \Pi(Q^2) = i \int d^4x e^{ip\cdot x} \langle 0 | T \{ \eta_N(x) \eta_N^\dagger(0) \} | 0 \rangle \tilde{d}_u \tilde{d}_d, \]

where \( Q^2 = -p^2 \), with \( p \) the current momentum.

We limit our calculation to the Lorentz structure proportional to \( \vec{p} \) because it is less susceptible to direct instanton contributions and excited resonances than the chirally even structure proportional to \( 1 \). Relevant diagrams for this correlator are shown in Fig. 2. After a straightforward calculation, we find

\[ \hat{p} \Pi^{\text{OPE}}(Q^2) = \frac{\hat{p}(\overline{q} q)}{\pi^2 f_\pi} \left[ \ln \left( \frac{\Lambda^2}{-p^2} \right) \pi^{\text{LO}} + \frac{1}{p^2} \ln \left( \frac{-p^2}{\Lambda_{\text{IR}}^2} \right) \pi^{\text{NLO}} + \frac{1}{p^2} \pi^{\text{NNLO}} \right]. \]

The leading order term is given by the diagrams 2a-2b,

\[ \pi^{\text{LO}} = \frac{3m_0^2}{64} \left[ d_- (5 - 2\beta - 3\beta^2) + d_+ (1 - \beta)^2 \right]. \]

Here we have introduced the combinations \( d_+ = \tilde{d}_u + \tilde{d}_d \) and \( d_- = \tilde{d}_u - \tilde{d}_d \). It turns out that the diagrams 2a, where the external source enters through the \( \langle 0 | \overline{q} \gamma \cdot \vec{p} q | 0 \rangle \) structure, give large and opposite sign contributions for the \( g_s \overline{q} \sigma q \) and \(-m_0^2 \overline{q} q \) sources so that their combined effect in \( H_q \) is nil. Fortunately, this cancelation does not hold for diagrams 2b that give (10).

The next-to-leading term corresponds to diagrams 2c,

\[ \pi^{\text{NLO}} = \frac{3d_-}{32} \left[ \frac{m_0^4 (1 - \beta^2) - g_s^2 (GG)}{9} (7 - 2\beta - 5\beta^2) \right] \]

\[ + \frac{d_+}{96} g_s^2 (GG) (1 - \beta)^2, \]

which contribute to the OPE (9) with the log of the infrared cutoff \( \Lambda_{\text{IR}} \).

\( g_s^2 (GG) \approx 0.4 - 1 \text{ GeV}^4 \) is the vacuum gluon condensate. The next-to-next-to-leading order

\[ \pi^{\text{NNLO}} = \frac{d_-}{24} \left[ \chi_S \pi^2 (1 + 2\beta - 3\beta^2) \right] \]

\[ + \frac{3m_0^4}{16} (13 - 2\beta - 11\beta^2) \right] + \frac{d_+}{24} \left[ \chi_T \pi^2 (1 - \beta)^2 \right] \]

\[ + \frac{m_0^4}{16} (-11 - 14\beta + \beta^2) - \frac{g_s^2 (GG)}{24} (5\beta^2 + 8\beta - 1) \].
contains the vacuum polarizabilities,
\[ \chi_{ST} \equiv \int d^4x \langle 0 | T \{ \pi u \pm \pi d(x), H_u \pm H_d(0) \} | 0 \rangle, \]  
and the vacuum factorization assumption has been made in (12).

The sum rules prescription involves matching the OPE with the phenomenological part, \( \Pi^{OPE}(Q^2) = \Pi^{\text{phen}}(Q^2) \), where
\[ \dot{p} \Pi^{\text{phen}} = \dot{p} \left( \frac{2\lambda^2 \bar{g}_{\pi NN} m_N}{(p^2 - m_N^2)^2} + \frac{A}{p^2 - m_N^2} \cdots \right) \]  
contains double and single pole contributions, and the continuum. After Borel transformation of the sum rule \( \Pi^{OPE}(Q^2) = \Pi^{\text{phen}}(Q^2) \) we obtain
\[ \frac{\langle \bar{q}q \rangle}{\pi^2 f_\pi} \left[ \pi^{\text{LO}} E_0 - \pi^{NLO} \left( \frac{M^2}{\Lambda^2_{IR}} - 0.58 \right) - \pi^{NNLO} \right] = \\
M^{-4} \exp \left[ -\frac{m_N^2}{M^2} \right] \left( 2\lambda^2 m_N \bar{g}_{\pi NN} + AM^2 \right) + M^{-2} B \exp \left[ -\frac{s}{M^2} \right] \]  
Here \( s \) is the continuum threshold and \( E_0 = 1 - e^{-s/M^2} \). \( A \) and \( B \) parametrize the contribution of excited states and are assumed to be independent of \( M \).

It is reasonable to start the numerical treatment from a simple estimate, à la Ioffe [21], which assumes the dominance of the ground state and the LO OPE term, and eliminates \( \lambda \) using the nucleon mass sum rule for \( \dot{p} \). Separating different isospin structures, we find
\[ \bar{g}_{\pi NN}^{(1)} = (\tilde{d}_u - \tilde{d}_d) \frac{3}{2} \frac{4\pi^2 \langle \bar{q}q \rangle |m_0^2|}{m_N f_\pi M^2} F_1(\beta) \]  
\[ \bar{g}_{\pi NN}^{(0)} = (\tilde{d}_u + \tilde{d}_d) \frac{3}{10} \frac{4\pi^2 \langle \bar{q}q \rangle |m_0^2|}{m_N f_\pi M^2} F_0(\beta). \]  
Here \( F_1(\beta) = (5 - 2\beta - 3\beta^2)/(5 + 2\beta + 5\beta^2) \) and \( F_0(\beta) = 5(1 - \beta)^2/(5 + 2\beta + 5\beta^2) \). \( F_1(0) = F_0(0) = 1 \). To get numerical estimates, we choose \( \beta = 0 \), extensively used in lattice simulations. It is well known that the \( j_1 \) current has a much better overlap with the nucleon ground state and \( \lambda_1 \gg \lambda_2 \).

Substituting \( M = 1\text{GeV} \), we obtain
\[ \bar{g}_{\pi NN}^{(1)} = 3 \times 10^{-12} \frac{\tilde{d}_u - \tilde{d}_d}{10^{-26}\text{cm} (225\text{MeV})^3} \frac{\langle \bar{q}q \rangle}{0.8\text{GeV}^2} \]  
\[ \bar{g}_{\pi NN}^{(0)} = 0.6 \times 10^{-12} \frac{\tilde{d}_u + \tilde{d}_d}{10^{-26}\text{cm} (225\text{MeV})^3} \frac{\langle \bar{q}q \rangle}{0.8\text{GeV}^2} \]  

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In most SUSY models, \( d_u(d) = \text{loop factor} \times M_{SUSY}^2 \times \text{a linear combination of } m_u \text{ and } m_d \). When combined with \( \langle \bar{q}q \rangle \) from eqs. (18-19), this forms \( m_\pi^2 f_\pi^2 \text{ times a function of } m_u/m_d \), thus eliminating a major source of uncertainty in EDM calculations due to the poor knowledge of \( m_u + m_d \) \([14, 15, 20]\). The estimate (18) is twice smaller than the value of \( \bar{g}^{(1)}_{\pi NN} \) used in \([3]\). Also in agreement with \([6]\), eqs. (18-19) suggest that \( \bar{g}^{(0)}_{\pi NN}/\bar{g}^{(1)}_{\pi NN} \sim 0.2 d_+ / d_- \).

For a more systematic analysis, one has to include NLO and NNLO terms in the OPE. Here, we immediately face the problem of the unknown vacuum condensates \( \chi_{S,T} \). Even though the vacuum correlators \( \langle \bar{q}q, \bar{q}q, G\sigma q \rangle \) can be determined using chiral perturbation theory \([22]\), there is no direct information on \( \langle \bar{q}q, \bar{q}g, G\sigma q \rangle \) other than that it is likely to be comparable with \( m_0^2 \langle \bar{q}q, \bar{q}q \rangle f_\pi^4 \). At this point we would like to take advantage of the possibility to choose \( \beta \) in such a way as to minimize higher order terms in the OPE. We note that \( \chi_S \) in (12) is multiplied by \( 1+2\beta-3\beta^2 \) which becomes 0 at \( \beta = -1/3 \) and 1. The choice of \( \beta = 1 \) also suppresses the leading order, while \( \beta = -1/3 \) maximizes it. For the expected size of \( \chi_S \) \([22]\), \( \chi_S \sim \pm 0.16 \times m_0^2 \langle \bar{q}q \rangle f_\pi^4 \), we can choose \( \beta \) in a range such that the whole square bracket in front of \( d_- \) in eq. (12) is zero. This gives a range of interpolating currents around \( \beta = -1/3 \),

\[-0.5 < \beta < 0, \tag{20}\]

where we can tune the NNLO terms to zero in the \( \bar{g}^{(1)}_{\pi NN} \) channel. Variation of \( \beta \) in this range contributes to an estimate of the uncertainty in our analysis. In the \( \bar{g}^{(0)}_{\pi NN} \) channel there is no obvious choice of \( \beta \) that would remove \( \chi_T \) and leave the leading order term un-suppressed, so we will choose the same \( \beta \) as for \( \bar{g}^{(1)}_{\pi NN} \). We note that this range is close to \( \beta = 0 \) as used for \([18,19]\).

One should also worry about the dependence on \( \Lambda_{IR} \) in NLO. Remarkably, in the range \([21]\), this dependence is softened by cancelation of the \( m_0^4 \) and \( g_s^2(GG) \) terms.

The preferred range for \( \bar{g}^{(1)}_{\pi NN} \) and \( \bar{g}^{(0)}_{\pi NN} \) is determined according to the following procedure. We take the OPE side of (13) at the lower point of the usual Borel window, \( M^2 = 0.8 \text{ GeV}^2 \), and vary it through the range of parameters \(-0.5 \leq \beta \leq 0, 300 \text{ MeV} \leq \Lambda_{IR} \leq 500 \text{ MeV}, 0.7 \text{ GeV}^2 \leq |m_0^2| \leq 0.9 \text{ GeV}^2, 0.4 \text{ GeV}^4 \leq g_s^2(GG) \leq 1 \text{ GeV}^4 \), and \( 2 \text{ GeV}^2 \leq s \leq 3 \text{ GeV}^2, 0.8 \text{ GeV}^6 \leq (2\pi)\lambda^2 \leq 0.9 \text{ GeV}^6 \) and finally \(-6 \text{ GeV}^{-1} \leq \chi_T / |\langle \bar{q}q \rangle| \leq 6 \text{ GeV}^{-1} \). On the r.h.s of (13) we assume the dominance of the double pole.
contributions for $M^2 = 0.8$ GeV$^2$ and allow for a 50% correction due to the presence of the unknown parameters $A$ and $B$, thus effectively widening the allowed range for $\bar{g}_{\pi NN}$. As expected, the couplings are most sensitive to the value of $m_0^2$. The final results are presented in Table 1. Our “best” value for $\bar{g}_{\pi NN}^{(1)}$ is determined by averaging over $\beta$ and choosing the central values for condensates, which also suppresses the logarithmic term. In order to separate the contribution of $\bar{g}_{\pi NN}^{(1)}$ from the $A$ and $B$ terms, we impose a relation among $A$, $B$ and $\pi^{LO}$ obtained by requiring the same large $M^2$ asymptotic behavior for both sides of (15). The resulting sum rule is fitted numerically and produces a result 1.5 times smaller than the naive estimate (18). For $\bar{g}_{\pi NN}^{(0)}$ the best value cannot be determined as the OPE side changes sign depending on the value of $\chi_T$.

Also included in this table are the preferred ranges for the CP-odd couplings of nucleons with $\eta$, $\rho$ and $\omega$ mesons. The couplings with $\rho$ and $\omega$ have the EDM-like structures $\frac{i}{2} \bar{N}(\partial_\nu V_\mu - \partial_\mu V_\nu)\sigma_{\mu\nu}\gamma_5 N$ with properly arranged isospin indices. They can be easily extracted from the calculation of the neutron EDM $d_n$, induced by $\tilde{d}_{u(d)}$ [15] after a simple re-assignment of charges for the external vector currents. Best values for $\tilde{g}_{\rho NN}$ and $\tilde{g}_{\omega NN}$ follow from the central values of $d_n(\tilde{d}_u,\tilde{d}_d)$ given in [15]. Finally, the coupling to the $\eta$ meson is dominated by the strange quark chromoelectric dipole moment $\tilde{d}_s$ in the isospin-singlet and by $\tilde{d}_u - \tilde{d}_d$ in the isospin-triplet channels, and in both cases only the expected range can be quoted.

In conclusion, we have shown that the size of the CP-odd pion-nucleon constant generated by quark chromoelectric dipole moments is given by the matrix element of $\bar{q}g_s G\sigma q - m_0^2 \bar{q}q$ over the nucleon state. We have constructed a QCD sum rule for this matrix element and determined the preferred range and the best value for the $\bar{g}_{\pi NN}^{(1)}$ coupling. The upper part of the preferred range agrees with previous estimates. However, in the interpretation of the experimental limit on the EDM of the mercury atom [3] in terms of limits imposed on new CP-violating physics, a more conservative value $\bar{g}_{\pi NN}^{(1)} = 2(\tilde{d}_u - \tilde{d}_d)/10^{-14}$ cm should be used. This translates the result of Ref. [8] (see [5, 9] for details) into the bound $|\tilde{d}_u - \tilde{d}_d| < 2 \times 10^{-26}$ cm. This constraint provide a sensitive probe of CP violation in the supersymmetric spectrum up to $M_{SUSY}$ of few TeV.

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Table 1: Preferred ranges and best values for the CP-odd meson-nucleon coupling constants induced by quark chromoelectric dipole moments. All values contain an overall multiplier $|\langle \bar{q}q \rangle|/(225 \text{ MeV})^3$. $d_{-,26}$, $d_{+,26}$ and $d_{s,26}$ are $d_-$ and $d_+$ and $d_s$ in units of $10^{-26} \text{ cm}$.

| coupling | preferred range | best value |
|----------|-----------------|------------|
| $g_{\pi NN}^{(1)} \times 10^{12}$ | (1 to 6) $d_{-,26}$ | 2 $d_{-,26}$ |
| $g_{\pi NN}^{(0)} \times 10^{12}$ | $(-0.5$ to $1.5)$ $d_{+,26}$ | -- |
| $g_{\eta NN}^{(0)} \times 10^{12}$ | $\sim (-1.5$ to $1.5)$ $d_{s,26}$ | -- |
| $g_{\eta NN}^{(1)} \times 10^{12}$ | $(-0.3$ to $1)$ $d_{-,26}$ | -- |
| $g_{\rho NN}^{(0)}$ | $(0.7$ to $2)$ $d_+$ | 1.4 $d_+$ |
| $g_{\rho NN}^{(1)}$ | $(0.5$ to $1.5)$ $d_-$ | 0.9 $d_-$ |
| $g_{\omega NN}^{(0)}$ | $(-1.5$ to $-0.5)$ $d_+$ | $-0.9$ $d_+$ |
| $g_{\omega NN}^{(1)}$ | $(-2$ to $-0.7)$ $d_+$ | $-1.4$ $d_+$ |

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