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Axial charges of nucleon resonances*

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Recently, first results have become available from lattice quantum chromodynamics (QCD) for two of the nucleon excitations, namely, the negative-parity $N^*(1535)$ and $N^*(1650)$ resonances [1]. The axial charge of the nucleon ground state had been studied before by different lattice-QCD groups in quenched calculations and with dynamical quarks [2–7]. In some of these works one has used chiral extrapolations (for a recent discussion of the associated problems see Ref. [8]), and the bulk of results obtained for g_A of the nucleon varies between about $1.10 \sim 1.40$.

Lately, the issue of axial constants of N^* resonances has become debated a lot due to the suggestion of chiral-symmetry restoration in the higher hadron spectra [9,10]. According to this scenario there should appear chiral doublets of positive- and negative-parity states and as a further consequence their axial charges should became small or almost vanishing. The first parity partners above the nucleon ground state are supposed to be the $N^*(1440)-N^*(1535)$, the next ones the $N^*(1710)-N^*(1650)$. The axial charges of the negative-parity partners in these pairs have been calculated in lattice QCD to be ~0.00 and ~0.55, respectively [1]; for the positive-parity states no results are yet available.

We have performed a study of the axial charges of N^* resonances in the framework of the relativistic constituent quark model (RCQM). Specifically we have extended a previous investigation of the nucleon axial form factors [11,12] to the first $J^P = \frac{1}{2}^\pm$ nucleon excitations. Our approach relies on solving the eigenvalue problem of the Poincaré-invariant mass operator in the framework of relativistic quantum mechanics. The axial current operator is chosen according to the spectator model (SM) [13]. For the RCQM we employed in the first instance the extended Goldstone-boson exchange (EGBE) RCQM [14], as it produces the most elaborate nucleon and N^* wave functions.

In Table 1 we present a selection of results for the axial charges g_A of the nucleon and the $N^*(1440)$, $N^*(1710)$, $N^*(1535)$, as well as $N^*(1650)$ resonances in case of the EGBE RCQM. It is immediately evident that the EGBE RCQM produces reasonable values for the axial charges in all instances without any further fittings. In the cases where a comparison is possible it produces the same pattern as lattice QCD. The g_A of the nucleon and of $N^*(1440)$ are practically of the same size, with the theoretical result for the nucleon being quite close to the experi-

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mental value of g_A=1.2695±0.0029 [15]. The nonrelativistic calculations cannot produce this value, neither in the simplistic $SU(6) \times O(3)$ quark model nor in the nonrelativistic limit of the RCQM. For the negative-parity $N^*(1535)$ resonance the g_A is predicted to be compatible with 0, while for the negative-parity $N^*(1650)$ resonance it is 0.51; both cases agree with the lattice-QCD results of Ref. [1]. Accidentally, the g_A value of the nonrelativistic $SU(6) \times O(3)$ quark model is similar in the $N^*(1650)$ case but the nonrelativistic limit of the EGBE RCQM shows deviations for both of the $\frac{1}{2}^-$ resonances. At this time nothing is known from lattice QCD for the $\frac{1}{2}^+$ resonances. For the latter, it would also be most interesting to check our results against lattice QCD, and we look forward to corresponding calculations.

Table 1. Predictions for axial charges g_A of the EGBE in comparison to available lattice QCD results [1-7], the values calculated by Glozman and Nefediev [9] within the $SU(6) \times O(3)$ nonrelativistic quark model, and the nonrelativistic limit from the EGBE RCQM.

State	J^P	EGBE	Lattice QCD	$SU(6) \times O(3)$ QM	EGBE nonrel
N(939)	$\frac{1}{2}^+$	1.15	1.10~1.40	1.66	1.65
N(1440)	$\frac{1}{2}^+$	1.16	–	1.66	1.61
N(1535)	$\frac{1}{2}^-$	0.02	~0.00	-0.11	-0.20
N(1710)	$\frac{1}{2}^+$	0.35	–	0.33	0.42
N(1650)	$\frac{1}{2}^-$	0.51	~0.55	0.55	0.64

It is particularly satisfying to find the RCQM predictions for the axial charges of the $N^*(1535)$ and $N^*(1650)$ resonances in agreement with the lattice-QCD results. We may thus be confident that at least for zero momentum-transfer processes the mass eigenstates of these nucleon excitations as produced especially with EGBE RCQM are quite reasonable. The latter is supposed to model the SBχS property of low-energy QCD. This type of hyperfine interaction, which also introduces an explicit flavor dependence, has been remarkably successful in describing a number of phenomena in low-energy baryon physics. Most prominently, it produces the correct level orderings of the positive- and negative-parity N^* resonances and simultaneously the ones in the other hyperon spectra, notably the Λ spectrum. The RCQM with GBE dynamics does not have any mechanism for chiral-symmetry restoration built in. As such it cannot be expected to produce parity doublets due to this reason. Nevertheless the EGBE RCQM describes the N^* resonance masses with good accuracy (mostly within the experimental error bars or at most exceeding them by 4%).

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Nucleon axial couplings and $[(\frac{1}{2},0) \oplus (0,\frac{1}{2})]$–$[(1,\frac{1}{2}) \oplus (\frac{1}{2},1)]$ chiral multiplet mixing*

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Abstract. Three-quark nucleon interpolating fields in QCD have well-defined, $U_A(1)$ and $SU_L(2) \times SU_R(2)$ chiral transformation properties: mixing of the $[(1,\frac{1}{2}) \oplus (\frac{1}{2},1)]$ chiral multiplet with one (of four available) $[(\frac{1}{2},0) \oplus (0,\frac{1}{2})]$, or $[(0,\frac{1}{2}) \oplus (\frac{1}{2},0)]$ fields can be used to fit the isovector axial coupling $g_A^{(1)}$ and thus predict the isoscalar axial coupling $g_A^{(0)}$ of the nucleon, in reasonable agreement with experiment. We also use a chiral meson-baryon interaction to calculate the masses and one-pion-interaction terms of $J = \frac{1}{2}$ baryons belonging to the $[(0,\frac{1}{2}) \oplus (\frac{1}{2},0)]$ and $[(1,\frac{1}{2}) \oplus (\frac{1}{2},1)]$ chiral multiplets and fit two of the diagonalized masses to the lowest-lying nucleon resonances thus predicting the third $J = \frac{1}{2}$ resonance at 2030 MeV, not far from the (one-star PDG) state $\Delta(2150)$.

1 Introduction

Almost 40 years ago Weinberg [1] considered mixing of chiral multiplets in the broken symmetry phase. In general such representation mixing may be complicated, but if only a few states are mixed, it may have predictive power. For instance, Weinberg used the mixing of $[(\frac{1}{2},0) \oplus (0,\frac{1}{2})]$ and $[(1,\frac{1}{2}) \oplus (\frac{1}{2},1)]$ to explain the nucleon's isovector axial coupling constant $g_A^{(1)} = 1.23$, its value at the time (the present value being 1.267). Weinberg's idea predated QCD and did not even invoke the existence of quarks, but it may still be viable in QCD. Indeed, this idea was revived in the early 1990's, since when it has been known by the name of mended symmetry [2].

The nucleon also has an isoscalar axial coupling $g_A^{(0)}$, which has been estimated from spin-polarized lepton-nucleon DIS data as $g_A^{(0)} = 0.28 \pm 0.16$ [3], or the more recent value $0.33 \pm 0.03 \pm 0.05$ [4]. The question is if the same chiral mixing angle can also explain the anomalously low value of this coupling? The answer manifestly depends on the $U_A(1)$ chiral transformation properties of the two admixed nucleon fields.

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In this paper we address this question using the $U_A(1)$ chiral transformation properties of nucleon fields [6,7] as derived from the three-quark nucleon interpolating fields in QCD. If the answer to our question turns out in the positive, we may speak about Weinberg's idea being viable in QCD. To test the present idea, besides the phenomenological study, we also investigate an extended linear sigma model containing baryon resonances, where we evaluate the axial couplings using baryon masses as input.

2 Three-quark nucleon interpolating fields

We start by summarizing the transformation properties of various quark trilinear forms with quantum numbers of the nucleon as shown in Refs. [6,7]. It turns out that every nucleon, i.e., spin- and isospin 1/2 field, besides having same non-Abelian transformation properties, comes in two varieties: one with "mirror" and another with "triple-naive" Abelian chiral properties.

In Table 1 we show the Abelian and non-Abelian chiral properties of the nucleon interpolating fields in QCD, Ref. [6,7]. Here we shall use those results as the theoretical input into our calculations. This constitutes a minimal assumption, as one has no other guide to the chiral representations of the nucleon. In Refs. [5–7] the local (non-derivative) spin $\frac{1}{2}$ baryon operators

$$N_1 = \epsilon_{abc}(\tilde{q}_a q_b)q_c, \tag{1}$$
$$N_2 = \epsilon_{abc}(\tilde{q}_a \gamma^5 q_b)\gamma^5 q_c, \tag{2}$$

were classified according to their Lorentz, chiral $SU_L(2) \times SU_R(2)$ (so-called "naive" chiral multiplet, whose axial charge is positive) and $U_A(1)$ group representations. Here we have introduced the "tilde-transposed" quark field $\tilde{q}$ as $\tilde{q} = q^T C\gamma_5(i\tau_2)$, where $C = i\gamma_2\gamma_0$ is the Dirac field charge conjugation operator, τ_2 is the second isospin Pauli matrix. Once one allows for the presence of one derivative, such as the so-called "mirror" $(0, \frac{1}{2}) \oplus (\frac{1}{2}, 0)$, whose axial charge is negative, Ref.[8],

$$N'_1 = \epsilon_{abc}(\tilde{q}_a q_b)i\partial_\mu\gamma^\mu q_c, \tag{3}$$
$$N'_2 = \epsilon_{abc}(\tilde{q}_a \gamma^5 q_b)i\partial_\mu\gamma^\mu\gamma^5 q_c. \tag{4}$$

and the $(1, \frac{1}{2}) \oplus (\frac{1}{2}, 1)$ nucleon chiral representation

$$N'_3 = i\partial_\mu(\tilde{q}\gamma_\nu q)\Gamma^{\mu\nu}_{3/2}\gamma_5 q, \tag{5}$$
$$N'_4 = i\partial_\mu(\tilde{q}\gamma_\nu\gamma_5\tau^i q)\Gamma^{\mu\nu}_{3/2}\tau^i q, \tag{6}$$

also become Pauli allowed, see Table 1. Here $\Gamma^{\mu\nu}_{3/2} = g^{\mu\nu} - \frac{1}{4}\gamma^\mu\gamma^\nu$. We found that indeed, as Gell-Mann and Levy [9] had postulated, the lowest-twist (non-derivative) J= $\frac{1}{2}$ nucleon field(s) form a $(\frac{1}{2}, 0)$ chiral multiplet, albeit there are two such independent fields. There is only one set of J= $\frac{1}{2}$ Pauli-allowed sub-leading-twist (one-derivative) interpolating fields that form a $(1, \frac{1}{2})$ chiral multiplet, however.

Table 1. The Abelian and the non-Abelian axial charges (+ sign indicates "naive", − sign "mirror" transformation properties) and the non-Abelian chiral mutiplets of $J^P = \frac{1}{2}$, Lorentz representation $(\frac{1}{2}, 0)$ nucleon fields. The field denoted by 0 belongs to the $(1, \frac{1}{2}) \oplus (\frac{1}{2}, 1)$ chiral multiplet and is the basic nucleon field that is mixed with various $(\frac{1}{2}, 0)$ nucleon fields in Eq. (7).

case	field	$g_A^{(0)}$	$g_A^{(1)}$	$SU_L(2) \times SU_R(2)$
I	$N_1 - N_2$	-1	$+1$	$(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$
II	$N_1 + N_2$	$+3$	$+1$	$(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$
III	$N_1' - N_2'$	$+1$	-1	$(0, \frac{1}{2}) \oplus (\frac{1}{2}, 0)$
IV	$N_1' + N_2'$	-3	-1	$(0, \frac{1}{2}) \oplus (\frac{1}{2}, 0)$
0	$N_3' + \frac{1}{3}N_4'$	$+1$	$+\frac{5}{3}$	$(1, \frac{1}{2}) \oplus (\frac{1}{2}, 1)$

3 Mixing of two chiral representations

Next consider the mixing of one of the fundamental chiral representations, as shown in Table 1 and the "higher" representation $(1, \frac{1}{2})$ for the nucleon,

$$\begin{aligned} g_{A\ \mathrm{mix.}}^{(1)} &= g_{A,\ \alpha}^{(1)} \cos^2\theta + g_{A\ (1,\frac{1}{2})}^{(1)} \sin^2\theta, \\ &= g_{A,\ \alpha}^{(1)} \cos^2\theta + \frac{5}{3} \sin^2\theta = 1.267. \end{aligned} \tag{7}$$

Here the suffix α corresponds to one of I-IV and the corresponding values of $g_{A,\ \alpha}^{(1)}$ are given in Table 1. We have also used the fact that $g_{A\ (1,\frac{1}{2})}^{(1)} = \frac{5}{3}$, see Ref. [1,7].

This provides a possible solution to the nucleon's axial coupling problem in QCD. Three-quark nucleon interpolating fields in QCD have well-defined two-fold $U_A(1)$ chiral transformation properties, see Table 1, that can be used to (naively) predict the isoscalar axial coupling $g_{A\ \mathrm{mix.}}^{(0)}$ as follows

$$g_{A\ \mathrm{mix.}}^{(0)} = g_{A,\ \alpha}^{(0)} \cos^2\theta + g_{A\ (1,\frac{1}{2})}^{(0)} \sin^2\theta, \tag{8}$$

together with the mixing angle θ extracted from Eq. (7). Note, however, that due to the different (bare) non-Abelian $g_A^{(1)}$ and Abelian $g_A^{(0)}$ axial couplings, see Table 1, the mixing formulae Eq. (8) give substantially different predictions from one case to another, see Table 2. We can see in Table 2 that the two candidates are cases I and IV, with $g_A^{(0)} = -0.2$ and $g_A^{(0)} = 0.4$, respectively, the latter being within 1-σ of the measured value $g_A^{(0)} = 0.33 \pm 0.08$. The nucleon field in case I is the well-known "Ioffe current", which reproduces the nucleon's properties in QCD lattice and sum rules calculations. The nucleon field in case IV is a "mirror" opposite of the orthogonal complement to the Ioffe current, an interpolating field that, to our knowledge, has not been used in QCD thus far.

3.1 A Simple Model

The next step is to try and reproduce this phenomenological mixing starting from a model interaction, rather than *per fiat*. As the first step in that direction we

Table 2. The values of the baryon isoscalar axial coupling constant predicted from the naive mixing and $g^{(1)}_{A\ \mathrm{expt.}} = 1.267$; compare with $g^{(0)}_{A\ \mathrm{expt.}} = 0.33 \pm 0.03 \pm 0.05$.

case	$(g^{(1)}_A, g^{(0)}_A)$	$g^{(1)}_{A\ \mathrm{mix.}}$	θ	$g^{(0)}_{A\ \mathrm{mix.}}$	$g^{(0)}_{A\ \mathrm{mix.}}$
I	$(+1, -1)$	$\frac{1}{3}(4 - \cos 2\theta)$	$\pm 39.3^o$	$-\cos 2\theta$	-0.20
II	$(+1, +3)$	$\frac{1}{3}(4 - \cos 2\theta)$	$\pm 39.3^o$	$2 + \cos 2\theta$	2.20
III	$(-1, +1)$	$\frac{1}{3}(1 - 4\cos 2\theta)$	$\pm 67.2^o$	1	1.00
IV	$(-1, -3)$	$\frac{1}{3}(1 - 4\cos 2\theta)$	$\pm 67.2^o$	$-(1 + 2\cos 2\theta)$	0.40

must look for a dynamical source of mixing. One such mechanism is the simplest chirally symmetric *non-derivative* one-(σ, π)-meson interaction Lagrangian, which induces baryon masses via its σ-meson coupling. We shall show that only the mirror fields couple to the $(1, \frac{1}{2})$ baryon chiral multiplet by non-derivative terms; the naive ones require one (or odd number of) derivative. This is interesting, as we have already pointed out that the mixing case IV seems a preferable one from the phenomenological consideration of axial couplings.

We use the projection method of Ref. [10] to construct the chirally invariant diagonal and off-diagonal meson-baryon-baryon interactions involving the "mirror" baryon $B_1 \in (0, \frac{1}{2})$, the $(B_2, \Delta) \in (1, \frac{1}{2})$ baryon and one $(\sigma, \pi) \in (\frac{1}{2}, \frac{1}{2})$ meson chiral multiplets. Here all baryons have spin 1/2, while the isospin of B_1 and B_2 is 1/2 and that of Δ is 3/2. The Δ field is then represented by an isovector-isospinor field Δ^i, $(i = 1, 2, 3)$. We found that for non-derivative mixing interaction the following chirally invariant combination

$$\mathcal{L}_3 = -g_3 \left[\bar{B}_1(\sigma + \frac{i}{3}\gamma_5 \tau \cdot \pi) B_2 + 4\bar{B}_1 i\gamma_5 \pi^i \Delta^i + \text{h.c.} \right], \qquad (9)$$

with the coupling constant g_3 induces an off-diagonal term in the baryon mass matrix after spontaneous symmetry breaking $\langle\sigma\rangle_0 \to f_\pi$ via its σ-meson coupling. Of course this is in addition to the conventional diagonal interactions:

$$\mathcal{L}_1 = -g_1 \bar{B}_1 \left(\sigma - i\gamma_5 \tau \cdot \pi\right) B_1, \qquad (10)$$

$$\mathcal{L}_2 = -\frac{2}{3} g_2 \Big[\bar{B}_2(\sigma + \frac{5}{3} i\gamma_5 \tau \cdot \pi) B_2 - 2\bar{\Delta}^i(\sigma + i\gamma_5 \tau \cdot \pi)\Delta^i - \frac{1}{\sqrt{3}} \bar{B}_2 \tau^i (\sigma + i\gamma_5 \tau \cdot \pi)\Delta^i + \text{h.c.} \Big], \qquad (11)$$

In writing down the Lagrangians (9,10,11), we have implicitly assumed that the parities of B_1, B_2 and Δ are the same. In principle, their parities are arbitrary, except for the parity of the ground state nucleon, which must be even. For instance, if B_2 has odd parity, the first term in the interaction Lagrangian Eq. (9) must include another γ_5 matrix. Here we consider all possible cases for the parities of B_2 and Δ.

Having established the mixing interaction Eq. (9), as well as the diagonal terms Eqs. (10),(11), we calculate the masses of the baryon states, as functions of the pion decay constant/chiral order parameter and (as yet undetermined)

Born approximation coupling constants. We diagonalize the mass matrix and express the mixing angle in terms of diagonalized masses. We find the following double-angle formulas for the mixing angles $\theta_{1,\dots,4}$, in the four different parities scenarios

$$\tan 2\theta_1 = \frac{\sqrt{(2N + \Delta)(2N^{*-} - \Delta)}}{(\Delta - N + N^{*-})}, \tag{12}$$

$$\tan 2\theta_2 = \frac{\sqrt{(\Delta - 2N)(2N^{*+} - \Delta)}}{(N + N^{*+} - \Delta)} \tag{13}$$

$$\tan 2\theta_3 = \frac{\sqrt{(2N - \Delta)(2N^{*-} + \Delta)}}{(\Delta - N + N^{*-})}, \tag{14}$$

$$\tan 2\theta_4 = \frac{\sqrt{-(\Delta + 2N)(2N^{*+} + \Delta)}}{(N + N^{*+} + \Delta)}, \tag{15}$$

where N is the nucleon ground state mass (940 MeV) and $N^{*\pm}, \Delta$ are the masses of the nucleon excited state, where $\pm$ indicates the parity of the N^* state. These angles correspond to the two (variable) parities as follows $\theta_1 \leftrightarrow (N^{*-}, \Delta^+)$, $\theta_2 \leftrightarrow (N^{*+}, \Delta^-)$, $\theta_3 \leftrightarrow (N^{*-}, \Delta^-)$, $\theta_4 \leftrightarrow (N^{*+}, \Delta^+)$, where $\pm$ indicates the parity of the state. Note that the angle θ_4 is necessarily imaginary so long as the Δ, N^* masses are physical (positive), and that the reality of the mixing angle(s) imposes stringent limits on the Δ, N^* resonance masses in other three cases, as well. Next, we use (some of) the observed resonance masses to determine the mixing angle(s) and thence calculate the axial couplings.

3.2 Results

Direct prediction

The four lowest-lying (besides the $N(940)$) candidate states in the PDG tables are: $R(1440)$, $N(1535)$, Δ(1620), Δ(1910), we use them to fit the free coupling constants. Of the two "mass allowed" scenarios, however, none survive the axial coupling test. Perhaps our choice of input resonances is inadequate. Note that one may "invert" this procedure, however, and use the isovector axial coupling to predict one of the baryon masses, say the Δ's, having fixed the other two, in this case the nucleon's $N(940)$ and $N^*(1440)$ or $N^*(1535)$.

Inverse prediction

Next, we use the double-angle formulas Eqs. (12)-(15) for the mixing angles $\theta_{1,\dots,4}$ together with the two observed nucleon masses to predict the Δ masses shown in the Table 3. We see that only the (N^{*+}, Δ^-) parity case leads to a realistic prediction: The difference between the observed (one-star) $S_{31}(2150)$ [11] Δ resonance mass and the predicted 2030 MeV may be neglected in view of the uncertainties and typical widths of states at such (high) energies. We shall not attach undue significance to this proximity in view of the rather uncertain status of this resonance, at least not until it is confirmed by another experiment. This choice of resonances leads to a reasonable πNN coupling constant (14.2 vs. 13.6 expt.) and predicts a set of as yet not measured π-baryon couplings.

Table 3. The values of the Δ baryon masses predicted from the isovector axial coupling $g^{(1)}_{A\ \text{mix.}} = g^{(1)}_{A\ \text{expt.}} = 1.267$ and $g^{(0)}_{A\ \text{mix.}} = 0.4$ vs. $g^{(0)}_{A\ \text{expt.}} = 0.33 \pm 0.08$.

$(N^{*P}, \Delta^{P'})$	$(N,\ N^*)$	Δ (MeV)	expt.
$(-,+)$	N(940), R(1535)	2330	1910
$(+,-)$	N(940), R(1440)	2030,2730	1620,2150
$(-,-)$	N(940), R(1535)	1140	1620,2150

A comment about the comparatively high value of the Δ mass seems to be in order now: In the mid-1960-s Hara [12] noticed that the chiral transformation rules for a $(1, \frac{1}{2})$ multiplet impose a strict and seemingly improbable mass relation among its two members: $m_\Delta = 2m_N$. The mixing with the $(\frac{1}{2}, 0)$ multiplet modifies this mass relation for the worse, i.e. it makes the Δ even heavier. For this reason, the lowest-lying Δ's of either parity cannot be the chiral partners of the nucleon ground state, as we initially assumed in our "direct prediction".

4 Three-field mixing

A linear superposition of yet another field (except for the mixture of cases II and III above) ought to give a perfect fit to both experimental values. Such an admixture introduces new free parameters (besides the two already introduced mixing angles, e.g. θ_1 and θ_4, we have the relative/mutual mixing angle θ_{14}, as the two nucleon fields I and IV may also mix). One may subsume/redefine the sum and the difference of the two angles θ_1 and θ_4 into the new angle θ, whereas one may define $\theta_{14} \doteq \varphi$ (this relationship depends on the precise definition of the mixing angles θ_1, θ_4 and θ_{14}); thus we find two equations with two unknowns of the general form:

$$\frac{5}{3}\sin^2\theta + \cos^2\theta \left(g_A^{(1)} \cos^2\varphi + g_A^{(1)\prime} \sin^2\varphi \right) = 1.267 \qquad (16)$$

$$\sin^2\theta + \cos^2\theta \left(g_A^{(0)} \cos^2\varphi + g_A^{(0)\prime} \sin^2\varphi \right) = 0.33 \qquad (17)$$

The values of the mixing angles obtained from this simple fit to the two baryon axial coupling constants are shown in Table 4. This, however, is not just a mere fit: when extending to the $SU_L(3) \times SU_R(3)$ symmetry, chiral transformation properties of the nucleon fields differ: $N_1 - N_2 \in [(\bar{3}, 3) \oplus (3, \bar{3})]$, $N_1 + N_2 \in [(8, 1) \oplus (1, 8)]$ and $(N_3' + \frac{1}{3}N_4') \in [(6, 3) \oplus (3, 6)]$, see Ref. [13]. From these chiral $SU_L(3) \times SU_R(3)$ symmetry assignment we can also predict the F and D couplings (un-corrected for the explicit $SU_F(3)$ symmetry breaking) in Table 4, which can be compared with the experimental numbers. We have not calculated the $SU_F(3)$ symmetry breaking corrections, as yet, so we have not taken into account the "error bars" on the mixing angle(s), which remains a task for the future. At any rate, it should be clear that the predicted values are "in the right ball park" for most of the scenarios considered here. Thus, the chiral multiplet mixing remains a viable theoretical scenario for the explanation of the nucleon isoscalar axial couplings.

Table 4. The values of the mixing angles obtained from the fit to the baryon axial coupling constants and the predicted values of axial F and D couplings. Experimental values have evolved from F=0.459 ± 0.008 and D=0.798 ± 0.008 in Ref. [14] to F=0.477 ± 0.001 and D=0.835 ± 0.001 in Ref. [15]. Note that the new values are more than 2-σ away from the old ones, and that the new F,D add up to F+D = 1.312 $\neq$ 1.269 ± 0.002.

case	$g_{A\ \text{expt.}}^{(0)}$	$g_{A\ \text{expt.}}^{(1)}$	θ	φ	F	D
I-II	1.267	0.33	39.3°	26.6°	0.399	0.868
I-III	1.267	0.33	49.6°	23.9°	0.333	0.934
I-IV	1.267	0.33	63.2°	53.9°	0.399	0.868

5 Summary and Discussion

We have shown that one can reproduce, within 1-σ uncertainty, the (unexpectedly small) isoscalar axial coupling of the nucleon by mixing (only) two (out of five independent) nucleon interpolating fields [1] by fitting the isovector- axial coupling. This solution to the nucleon spin problem does not invoke exotica such as a) hidden strangeness; or b) polarized gluon components in the nucleon wave function, in agreement with recent results of the COMPASS experiment [16],[17]. This scenario is quantitatively reproduced in a simple dynamical model which then predicts the existence of the S_{31} resonance at 2160 MeV, in agreement with the PDG tables [11]. By mixing three nucleon interpolating field chiral multiplets one may simultaneously fit both the isovector and the isoscalar axial couplings and predict the SU(3) F and D couplings, which have the correct size within the expected $\mathcal{O}(20\%)$ SU(3) symmetry breaking corrections.

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