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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Three-quark nucleon interpolating fields in QCD have well-defined $SU_L(2)\times SU_R(2)$ and $U_A(1)$ chiral transformation properties. Mixing of the $[(1,\frac{1}{2})\oplus(\frac{1}{2},1)]$ chiral multiplet with one of $[(\frac{1}{2},0)\oplus(0,\frac{1}{2})]$ or $[(0,\frac{1}{2})\oplus(\frac{1}{2},0)]$ representation 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 $\Delta(2150)$.

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Introduction Chiral symmetry, as one of the symmetries of QCD, is a key property of the strong interactions. When the chiral $SU_L(2)\times SU_R(2)$ symmetry is spontaneously broken to $SU_V(2)$, the broken chiral symmetry plays a dynamical role in the low energy theorems among the Nambu-Goldstone bosons, i.e. the pions. Hadrons are then classified according to the residual symmetry $SU_V(2)$.

Almost 40 years ago Weinberg [1] proposed a scenario where the consideration of the full chiral symmetry group makes sense in the broken symmetry phase. There hadron states are represented by the representations of the chiral symmetry group but with representation mixing. In general such mixing is complicated in the broken symmetry phase, but if it can be described by a few parameters, it may have predictive power. For instance, the nucleon's isovector axial coupling constant is determined by its chiral representation [1, 2]. Weinberg then considered the mixing of $[(\frac{1}{2},0)\oplus(0,\frac{1}{2})]$ and $[(1,\frac{1}{2})\oplus(\frac{1}{2},1)]$ as an explanation of the nucleon's isovector axial coupling constant $g_A^{(1)}=1.23$, its value at the time (the present value being 1.267) [18]. Once various chiral representations are included, they also have relevance to the physics of excited states as well as the ground state. 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]. Related development was also made where a particular representation so called mirror representation $[(0,\frac{1}{2})\oplus(\frac{1}{2},0)]$ was introduced and physical relevance was discussed [3, 4].

The nucleon also has an isoscalar axial coupling $g_A^{(0)}$, which has been estimated from spin-polarized lepton-nucleon DIS data as $g_A^{(0)}=0.28\pm0.16$ [5], or the more recent value $0.33\pm0.03\pm0.05$ [8]. The question is if the same chiral mixing angles can also explain the anomalously low value of this coupling? The answer manifestly depends on the $U_A(1)$ chiral transformation properties of the admixed nucleon fields.

In this Letter we address these question about axial couplings and some other properties of baryon excited states using the $SU_L(2)\times SU_R(2)$ and $U_A(1)$ chiral transformation properties of nucleon interpolating fields [6, 7] as derived from the three-quark nucleon interpolating fields in QCD. Here we use the properties of the nucleon fields as a guide for those of the corresponding states. 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, we also investigate an extended linear sigma model containing baryon resonances, where we evaluate the axial couplings using baryon masses as input.

Basic facts and assumptions Let us start our discussion by writing down the following three-quark nucleon interpolating fields:

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

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. Properties of these particular forms were investigated in Refs. [6, 7]. The local (non-derivative) spin $\frac{1}{2}$ baryon operators were clas-

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TABLE I: The Abelian and the non-Abelian axial couplings (+ 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	$N_3^{'} + \frac{1}{3} N_4^{'}$	$+1$	$+\frac{5}{3}$	$(1, \frac{1}{2}) \oplus (\frac{1}{2}, 1)$

sified according to their Lorentz, chiral $SU_L(2) \times SU_R(2)$ and $U_A(1)$ group representations. The chiral representation of Eqs. (1,2) are both the so-called "naive", $[(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})]$ and characterized by the positive axial coupling constant. In the present fundamental representation it is unity. Properties of the Abelian ($U_A(1)$) and non-Abelian ($SU_L(2) \times SU_R(2)$) chiral symmetries are summarized in Table I, 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.

If 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 coupling is negative, Ref.[4] [19],

$$N_1' = \epsilon_{abc}(\tilde{q}_a q_b) i\partial_\mu \gamma^\mu q_c, \quad (3)$$
$$N_2' = \epsilon_{abc}(\tilde{q}_a \gamma^5 q_b) i\partial_\mu \gamma^\mu \gamma^5 q_c, \quad (4)$$

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

$$N_3' = i\partial_\mu(\tilde{q}\gamma_\nu q)\Gamma^{\mu\nu}_{3/2}\gamma_5 q, \quad (5)$$
$$N_4' = i\partial_\mu(\tilde{q}\gamma_\nu\gamma_5\tau^i q)\Gamma^{\mu\nu}_{3/2}\tau^i q, \quad (6)$$

also become available, see Table I. 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. [20] 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. Here we note that the mirror and higher representations can also be made of multiquark (five or more) fields. The consideration of such multiquark components is also interesting, but lies beyond the scope of the present study.

Let us now consider the mixing of one of the fundamental chiral representations, as shown in Table I, and the "higher" representation $(1, \frac{1}{2})$ for the nucleon. The diagonal component of the axial coupling constant in the mixed state is calculated as follows [1]

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

Here the suffix α corresponds to one of I-IV and the corresponding values of $g^{(1)}_{A,\ \alpha}$ are given in Table I. We have also used the fact that $g^{(1)}_{A\ (1,\frac{1}{2})} = \frac{5}{3}$, see Ref. [1, 7]. Several comments are called for now: 1) a tacit assumption underlying Eq. (7) is that the axial coupling(s) of the baryon fields do not change due to the shift from the Wigner-Weyl to the Nambu-Goldstone phase (and *vice versa*); 2) assumption no. 1) is related to the assumption that no part of the axial current is induced by derivative interactions of the Bjorken-Nauenberg type [10]; 3) although assumption no. 1) need not be a good one, we nevertheless retain it, following Weinberg [1], as relaxing it would require new free parameters to be introduced.

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 I, that can be used to predict the isoscalar axial coupling $g^{(0)}_{A\ \mathrm{mix.}}$ as follows

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

together with the mixing angle θ extracted from Eq. (7). Note, however, that due to the different non-Abelian $g_A^{(1)}$ and Abelian $g_A^{(0)}$ axial couplings, the mixing formula Eq. (8) give substantially different predictions from one case to another, see Table II. We can see in Table II 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.

TABLE II: 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	$(g_A^{(1)}, g_A^{(0)})$	$g^{(1)}_{A\ \mathrm{mix.}}$	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

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 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 with the mirror baryon seems a preferable one from the phenomenological consideration of axial couplings.

We use the projection method of Ref. [11] 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 $SU_L(2) \times SU_R(2)$ chirally invariant combination

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

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 [11]:

$$\begin{aligned} \mathcal{L}_1 &= -g_1 \bar{B}_1 \left(\sigma - i \gamma_5 \tau \cdot \pi \right) B_1, & (10) \\ \mathcal{L}_2 &= -\frac{2}{3} g_2 \Big[\bar{B}_2 (\sigma + \frac{5}{3} i \gamma_5 \tau \cdot \pi) B_2 \\ &\quad -2 \bar{\Delta}^i (\sigma + i \gamma_5 \tau \cdot \pi) \Delta^i \\ &\quad - \frac{1}{\sqrt{3}} \bar{B}_2 \tau^i (\sigma + i \gamma_5 \tau \cdot \pi) \Delta^i + h.c. \Big] . & (11) \end{aligned}$$

In the last term of (11), the $\sigma N \Delta$ interaction vanishes, but that form preserves chiral invariance. 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, they are arbitrary, except for 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 [4]. 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) and (11), we calculate the masses of the baryon states, as functions of the pion decay constant/chiral order parameter and the coupling constants g_1, g_2 and g_3. 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 $\theta_{1,\ldots,4}$ between B_1 and B_2 in the four different parities scenarios

$$\begin{aligned} \tan 2\theta_1 &= \frac{\sqrt{(2N + \Delta)(2N^* - \Delta)}}{(\Delta - N + N^*)}, & (12) \\ \tan 2\theta_2 &= \frac{\sqrt{(\Delta - 2N)(2N^* - \Delta)}}{(N + N^* - \Delta)}, & (13) \\ \tan 2\theta_3 &= \frac{\sqrt{(2N - \Delta)(2N^* + \Delta)}}{(\Delta - N + N^*)}, & (14) \\ \tan 2\theta_4 &= \frac{\sqrt{-(\Delta + 2N)(2N^* + \Delta)}}{(N + N^* + \Delta)}, & (15) \end{aligned}$$

where N, N^* and Δ represent the masses of the corresponding particles. The four angles correspond to the four possibel parities; $\theta_1 : (N^{*-}, \Delta^+)$, $\theta_2 : (N^{*+}, \Delta^-)$, $\theta_3 : (N^{*-}, \Delta^-)$ and $\theta_4 : (N^{*+}, \Delta^+)$, where $\pm$ indicate 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.

In the present study we have three model parameters g_1, g_2 and g_3, which can be determined by different set of inputs. In the following we consider two cases. The first case uses three baryon masses as inputs (Direct prediction) and determine the mixing angles for the prediction of the axial couplings. The second case uses two baryon masses and the mixing angle as inputs and predicts the third baryon mass (Inverse prediction).

TABLE III: The particle assignments and parities. The values in brackets are from PDG [12].

	Assignment	Mass(Exp)
Δ^{+i}	P_{31}	(1910)
Δ^{-i}	S_{31}	(1620)
N^{*-}	S_{11}	(1535)
N^{*+}	P_{11}	(1440)

Direct prediction The four lowest-lying (besides the $N(940)$) candidate states in the PDG tables are: $N^*(1440)$, $N^*(1535)$, $\Delta(1620)$ and $\Delta(1910)$ (Table III). We use them to fit the free coupling constants. Only two out of four scenarios (1 and 3) turn out to be viable, at least as far as the masses are concerned; the second and fourth scenarios predict imaginary coupling constants, i.e., non-Hermitian coupling "Hamiltonian", due to baryon masses that do not obey the constraints of chiral symmetry, see Table IV. Of the two allowed scenarios, however, none survive the axial coupling test as shown in the last three columns in Table IV, if we use the g_A values as listed in Table I. In this case our choice of input resonances, Table III is inadequate. Note that the Lagrangians (9-11) are more general than expected from the structure of the three-quark fields.

Therefore, we "invert" this procedure 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 $N^*(1535)$.

TABLE IV: The values of the free parameters (theoretical coupling constants) and the mixing angle θ extracted from the baryon masses in various scenarios. Note that only the absolute values of g_3 and θ can be extracted from this analysis. In the last two columns we show the axial couplings $g_A^{(1)}$ and $g_A^{(0)}$, for the ground-state nucleon field with bare (unmixed) axial coupling $g_A^{(0)}$=-3.

$(N^{*P}, \Delta^{iP'})$	g_1	g_2	$\pm g_3$	$\pm\theta$	$g_A^{(1)}(N)$	$g_A^{(0)}(N)$
$(-,+)$	-3.87	15.4	5.64	28.9^o	-0.38	-2.06
$(+,-)$	16.9	13.1	1.54 i	$28.1i^o$	-1.69	-4.04
$(-,-)$	-2.31	-13.1	6.06	32.8^o	-0.22	-1.83

Inverse prediction Next, we use the formulas Eqs. (12)-(15) for the (double) mixing angles $\theta_{1,...,4}$ together with the two observed nucleon masses to predict the Δ masses shown in the Table V. We see that only the (N^{*+}, Δ^-) parity case leads to a realistic prediction. The difference between the observed (one-star) $S_{31}(2150)$ [12] Δ resonance mass and the predicted 2030 and 2730 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 mixing angle automatically leads to a reasonable πNN coupling constant (12.8 vs. 13.6 expt.), due to the validity of the Goldberger-Treiman relation, but also predicts a set of as yet not measured π-baryon couplings, see Table VI, that can be used to test the model.

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

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

TABLE VI: The values of the physical π-baryon and axial coupling constants predicted in the only physically viable scenario II(+,-), i.e. with $N^* = R(1440)$ and Δ(2150). Here $g^{(1)}_{A\ \mathrm{expt.}} = 1.267$ and $g^{(0)}_{A\ \mathrm{mix.}} = 0.4$.

(P, P')	$g_{\pi NN}$	$g_{\pi RR}$	$g_{\pi NR}$	$g_{\pi N\Delta}$	$g_{\pi R\Delta}$	$g_{\pi\Delta\Delta}$
$(+,-)$	12.8	-9.3	-12.2	1.10	7.29	-21.8

A comment about the comparatively high value of the Δ mass seems to be in order now: In the mid-1960-s Hara [13] 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 cannot be a chiral partner of the lowest-lying nucleon field, whereas, in Δ(2150) we seem to have found a reasonable candidate for the $N(940)$'s chiral partner.

Three-field mixing In a more realistic analysis, we may consider mixings of various independent fields, for example, five fields as shown in Table I. This, however, means that there are $5 \times 4/2 = 10$ angles that parametrize a 5×5 real, orthogonal O(5) matrix. Manifestly, such a proliferation of free parameters allows a much greater freedom in fitting the data, but also allows appearance of ambiguities with the present day paucity of data. For this reason we shall confine ourselves to mixing (only) one component at a time.

Manifestly, a linear superposition of yet another field (except for the mixture of cases II and III in Table I) ought to give a fit to both experimental values. Such an admixture of the three fields of chiral representations I, IV and 0 of Table I 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 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:

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

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

The values of the mixing angles obtained from this simple fit to the two baryon axial coupling constants are shown in Table VII. 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 in the values of their (bare) F and D coefficients: $N_1 - N_2 \in [(\bar{3}, 3) \oplus (3, \bar{3})]$ has $D = 1, N = 0$, $N_1 + N_2 \in [(8,1) \oplus (1,8)]$ has $D = 0, N = 1$, and $(N_3' + \frac{1}{3}N_4') \in [(6,3) \oplus (3,6)]$ has $D = 1, N = \frac{2}{3}$ [14]. From these chiral $SU_L(3) \times SU_R(3)$ symmetry assignments we can independently "predict" the physical (mixed) F and D couplings in Table VII, which can be compared with the experimental values. We have not calculated the $SU_F(3)$ symmetry breaking corrections, as yet, so we could not take 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". Thus, the chiral multiplet mixing remains a viable theoretical scenario for the explanation of the nucleon properties including isoscalar axial couplings.

Summary, Comments We have shown that one can

TABLE VII: 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. [15].

case	$g^{(1)}_{A\ \mathrm{expt.}}$	$g^{(0)}_{A\ \mathrm{expt.}}$	θ	φ	F	D
I-II	1.267	0.33	39.3^o	26.6^o	0.399	0.868
I-III	1.267	0.33	50.7^o	23.9^o	0.399	0.868
I-IV	1.267	0.33	63.2^o	53.9^o	0.399	0.868

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 [21] by fitting the isovector- axial coupling. This solution to the nucleon spin problem does not invoke exotica such as a) polarized strange sea quarks; 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 [12]. 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.

This work was inspired by Weinberg's early work, from which it differs in several significant ways: 1) the admixed nucleon fields have been explicitly written out in terms of three-quark interpolators of QCD; 2) their chiral properties were extracted from the interpolators; 3) as a consequence of point 2) some unexpected $\mathrm{U}_A(1)$ and SU(3) properties were obtained and then used to calculate the flavor-singlet and the F and D flavor-octet axial couplings; 4) the mixing angle as a function of the baryon masses is based on a simple chiral interaction of baryons and spinless mesons; 5) only $J = \frac{1}{2}$ Δ and nucleon resonances were used.

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[18] In the notation $[(I_L, I_R) \oplus (I_R, I_L)]$, the first (I_L, I_R) denote the chiral multiplet $SU_L(2) \times SU_R(2)$ for the left-handed nucleon and the second one for the right-handed nucleon.

[19] For color gauge invariance, we need the covariant derivative. However, the present discussion of chiral properties is not modified by the covariant derivative.

[20] In what follows we indicate the chiral representation only for the left-handed component.

[21] or equivalently relativistic component in the nucleon's Bethe-Salpeter wave function