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Nucleon magnetic moments in an extended chiral constituent quark model

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Abstract. We present results for the nucleon magnetic moments in the context of an extended chiral constituent quark model based on the mechanism of the Goldstone boson exchange, as suggested by the spontaneous breaking of chiral symmetry in QCD. The electromagnetic charge-current operator is consistently deduced from the model Hamiltonian, which includes all force components for the pseudoscalar, vector and scalar meson exchanges. Thus, the continuity equation is satisfied for each piece of the interaction, avoiding the introduction of any further parameter. A good agreement with experimental values is found. The role of isoscalar two-body operators, not constrained by the continuity equation, is also investigated.

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The main problem for the realization of QCD in the non-perturbative sector of strong interactions is the identification of the effective degrees of freedom that pertain to the low-energy properties of hadrons. By considering the features of the baryon spectrum, the well-known phenomenon of spontaneous breaking of the chiral symmetry in QCD suggests that below a certain energy threshold a baryon should be represented in terms of two types of effective "particles": the constituent quark, whose dynamical mass is related to the $\langle q\bar{q}\rangle$ condensate, and the Goldstone bosons, which couple directly to the constituent quarks [1].

The mechanism of Goldstone boson exchange has successfully been implemented in a constituent quark model [2,3] (for the sake of brevity, indicated as "the model"). The peculiar spin-flavor dependence of the hyperfine interaction correctly reproduces the ordering of the positive- and negative-parity bands of light and strange baryons in a unified manner; a long-standing problem of baryon spectroscopy that no other constituent quark model (either based on the traditional gluon exchange mechanism [4–12], or with the additional inclusion of contributions from meson exchanges [13–16]) has ever been able to solve. The Hamiltonian of the model is given by

$$H = \sqrt{\mathcal{M}^2 + \mathbf{P}^2},$$
$$\mathcal{M} = \sum_{i=1}^{3} \sqrt{\mathbf{y}_i^2 + m_i^2} + \sum_{i<j=1}^{3} V_{ij}, \qquad (1)$$

where $\mathbf{P} = \sum_{i=1}^{3} \mathbf{p}_i$ is the center-of-mass (cm) momentum of the three constituent quarks with mass m_i and momentum $\mathbf{p}_i$. The mass operator $\mathcal{M}$ describes the intrinsic motion of the quarks inside the baryon in terms of their intrinsic momenta $\mathbf{y}_i = \mathbf{p}_i - (1/3)\mathbf{P}$ and mutual interaction V_{ij}, where the (multi) Goldstone boson exchange is realized by including all the force components of the pseudoscalar (π, K, η, η'), vector $(\rho, K^*, \omega_8, \omega_0)$ and scalar (σ) meson exchanges. The use of the relativistic expression for the kinetic energy operator avoids the typical drawback of the nonrelativistic constituent quark models, where the mean velocity of quarks can become larger than the velocity of light. Once m_i are fixed to 340 MeV, the model has five fitting parameters that are determined by the observed nucleon mass and baryon spectra.

A first test of the model wave function is to calculate the nucleon rms radii. However, according to common practice the nucleon radius is deduced from the slope at the origin of the Q-dependent charge form factor. This would necessarily involve additional parameters to describe the Q-dependence of the internal structure of the effective degrees of freedom represented by the constituent

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$$\mathbf{J}_{\mathrm{intr}}^{\mathrm{drift}} = \int \mathrm{d}\mathbf{x}\, \mathrm{e}^{\mathbf{i}\mathbf{q}\cdot\mathbf{x}} \langle f| \sum_{l=1}^{3} \frac{\delta}{\delta \mathbf{A}_l(\mathbf{x})} \sum_{i=1;\, j,k\neq i}^{3} \sqrt{\left[\frac{2}{3}\big(\mathbf{p}_i - e_i\mathbf{A}_i(\mathbf{x})\big) - \frac{1}{3}\big(\mathbf{p}_j - e_j\mathbf{A}_j(\mathbf{x})\big) - \frac{1}{3}\big(\mathbf{p}_k - e_k\mathbf{A}_k(\mathbf{x})\big)\right]^2 + m_i^2}\,\Bigg|_{\mathbf{A}_l=0} |i\rangle$$

$$= \langle f| \sum_{i=1;\, j,k\neq i}^{3} \frac{\mathbf{y}_i + \mathbf{y}'_i}{E_i + E'_i} \left[\frac{2}{3} e_i \delta\big(\mathbf{p}'_i - \mathbf{p}_i - \mathbf{q}\big)\delta\big(\mathbf{p}'_j - \mathbf{p}_j\big)\delta\big(\mathbf{p}'_k - \mathbf{p}_k\big) - \frac{1}{3} e_j \delta\big(\mathbf{p}'_j - \mathbf{p}_j - \mathbf{q}\big)\delta\big(\mathbf{p}'_k - \mathbf{p}_k\big)\delta\big(\mathbf{p}'_i - \mathbf{p}_i\big)\right.$$

$$\left. - \frac{1}{3} e_k \delta\big(\mathbf{p}'_k - \mathbf{p}_k - \mathbf{q}\big)\delta\big(\mathbf{p}'_i - \mathbf{p}_i\big)\delta\big(\mathbf{p}'_j - \mathbf{p}_j\big)\right] |i\rangle, \tag{5}$$

quarks. In a forthcoming paper we will show that the model gives reasonable values provided that form factors are introduced for the constituent quarks too.

As a further test we will consider here the magnetic moments of the nucleon. In the naive nonrelativistic constituent quark model the octet baryon magnetic moments are obtained by summing up the Dirac magnetic moments of three free quarks with results as close as 5–10% to the experimental data [17]. Many theoretical attempts to improve the approach involve configuration mixing, relativistic corrections, isospin-violating effects [18–21] as well as pionic effects [22–24]. In all these calculations a better agreement of the results with data is obtained at the price of introducing further parameters and/or assuming simplifying assumptions on the baryon wave function or the current operator. Relativistic constituent quark models have also been studied using the light-front formalism [10, 25–27]. Relativistic effects can be of order 20%. However, in all these calculations it was impossible to fit simultaneously the proton and neutron magnetic moments without some sort of modification of the quark model parameters and/or introducing anomalous quark moments. In the present letter the nucleon magnetic moments are obtained by calculating the matrix elements of the electromagnetic charge-current operator deduced consistently from the Hamiltonian H of eq. (1). In this way, the charge-current operator is gauge invariant, satisfies the continuity equation and no further parameters are introduced. The initial and final states ($|i\rangle, |f\rangle$) are taken as the factorized product of the eigenfunctions of $\mathcal{M}$ and of plane waves for the cm motion.

The electromagnetic charge-current operator consists of a one- and a two-body part. Since the kinetic energy operator of the cm and intrinsic motion in eq. (1) contains the square root operator and the potential is local, a gauge invariant one-body operator is deduced by applying the minimal substitution with an external electromagnetic field $A = (A_0, \mathbf{A})$ and then by using the formalism of the functional derivation. In fact, using the results of ref. [28], we can define

$$H(A) = \sqrt{\mathcal{M}^2(\mathbf{A}) + \left[\sum_i \big(\mathbf{p}_i - e_i\mathbf{A}\big)\right]^2} + e_{\mathrm{N}} A_0$$
$$\equiv \sqrt{\mathcal{M}^2(\mathbf{A}) + R^2(\mathbf{A})} + e_{\mathrm{N}} A_0, \tag{2}$$

where e_i are the individual quark charges and e_{N} is the nucleon charge (both expressed in units of the proton charge), and then represent the matrix element of the one-body charge-current operator as

$$J_0 = \int \mathrm{d}\mathbf{x}\, \mathrm{e}^{\mathbf{i}\mathbf{q}\cdot\mathbf{x}} \langle f| \frac{\delta H(A)}{\delta A_0(\mathbf{x})}\bigg|_{A_0=0} |i\rangle$$
$$= e_{\mathrm{N}} \delta\big(\mathbf{P}' - \mathbf{P} - \mathbf{q}\big), \tag{3}$$

$$\mathbf{J}^{\mathrm{drift}} = \int \mathrm{d}\mathbf{x}\, \mathrm{e}^{\mathbf{i}\mathbf{q}\cdot\mathbf{x}} \langle f| \frac{\delta H(A)}{\delta \mathbf{A}(\mathbf{x})}\bigg|_{\mathbf{A}=0} |i\rangle$$
$$= \frac{1}{E+E'} \int \mathrm{d}\mathbf{x}\, \mathrm{e}^{\mathbf{i}\mathbf{q}\cdot\mathbf{x}} \langle f| \frac{\delta H^2(A)}{\delta \mathbf{A}(\mathbf{x})}\bigg|_{\mathbf{A}=0} |i\rangle$$
$$= \frac{2M}{E+E'} \int \mathrm{d}\mathbf{x}\, \mathrm{e}^{\mathbf{i}\mathbf{q}\cdot\mathbf{x}} \langle f| \frac{\delta \mathcal{M}(A)}{\delta \mathbf{A}(\mathbf{x})}\bigg|_{\mathbf{A}=0} |i\rangle$$
$$+ e_{\mathrm{N}} \frac{\mathbf{P} + \mathbf{P}'}{E+E'}\, \delta\big(\mathbf{P}' - \mathbf{P} - \mathbf{q}\big)$$
$$\equiv \frac{2M}{E+E'}\, \mathbf{J}_{\mathrm{intr}}^{\mathrm{drift}} + e_{\mathrm{N}} \frac{\mathbf{P} + \mathbf{P}'}{E+E'}\, \delta\big(\mathbf{P}' - \mathbf{P} - \mathbf{q}\big). \tag{4}$$

Here, $\mathbf{q}$ is the momentum transferred by the external field at the space point $\mathbf{x}$, M is the nucleon mass and $\mathbf{P}(\mathbf{P}')$, $E(E')$ are the cm momentum and total energy of the initial (final) state $|i(f)\rangle$, respectively. The spatial part (4) represents the contribution of the total drift current: it contains a cm part, that describes the nucleon as a whole, and a part $\mathbf{J}_{\mathrm{intr}}^{\mathrm{drift}}$ related to the intrinsic motion. The latter can be made explicit by again systematically applying the minimal substitution to each quark momentum variable and then using the techniques of functional derivation [28]:

see equation (5) *above,*

where $E_i = \sqrt{\mathbf{y}_i^2 + m_i^2}$ and $\mathbf{y}'_i = \mathbf{y}_i + (2/3)\mathbf{q}$.

Following the lines of ref. [28], the one-body spin magnetic current can also be deduced by applying the minimal substitution to the equivalent Hamiltonian

$$H = \sqrt{\mathcal{M}_s^2 + (\boldsymbol{\sigma}\cdot\mathbf{P})^2} \equiv \sqrt{\mathcal{M}_s^2 + R_s^2},$$
$$\mathcal{M}_s = \sum_{i=1}^{3} \sqrt{(\boldsymbol{\sigma}\cdot\mathbf{y}_i)^2 + m_i^2} + \sum_{i<j=1}^{3} V_{ij} \tag{6}$$

and by defining

$$\begin{aligned}\mathbf{J}^{\rm spin} &= \mathbf{J}^{\rm tot} - \mathbf{J}^{\rm drift} \\ &= \int \mathrm{d}\mathbf{x}\, \mathrm{e}^{\mathrm{i}\mathbf{q}\cdot\mathbf{x}} \langle f| \frac{\delta}{\delta \mathbf{A}(\mathbf{x})} \Bigg[\sqrt{\mathcal{M}_s^2(A) + R_s^2(A)} \\ &\quad - \sqrt{\mathcal{M}^2(A) + R^2(A)} \Bigg] \Bigg|_{\mathbf{A}=0} |i\rangle \\ &= \int \mathrm{d}\mathbf{x}\, \mathrm{e}^{\mathrm{i}\mathbf{q}\cdot\mathbf{x}} \Bigg\{ \frac{1}{E+E'} \langle f| \frac{\delta}{\delta \mathbf{A}(\mathbf{x})} \\ &\quad \times \big[R_s^2(A) - R^2(A)\big] \Big|_{\mathbf{A}=0} |i\rangle + \frac{2M}{E+E'} \langle f| \frac{\delta}{\delta \mathbf{A}(\mathbf{x})} \\ &\quad \times \big[\mathcal{M}_s(A) - \mathcal{M}(A)\big] \Big|_{\mathbf{A}=0} |i\rangle \Bigg\} \\ &\equiv \mathbf{J}^{\rm spin}_{\rm cm} + \mathbf{J}^{\rm spin}_{\rm intr}. \end{aligned} \tag{7}$$

After some algebra, the final result for the cm and intrinsic one-body spin currents is

$$\begin{aligned}\mathbf{J}^{\rm spin}_{\rm cm} &= \frac{\mathrm{i}}{E+E'} \langle f| \sum_{i=1;\, j,k\neq i}^{3} e_i \boldsymbol{\sigma}_i \times (\mathbf{p}'_i - \mathbf{p}_i) \\ &\quad \times \delta(\mathbf{p}'_i - \mathbf{p}_i - \mathbf{q}) \delta(\mathbf{p}'_j - \mathbf{p}_j)\, \delta(\mathbf{p}'_k - \mathbf{p}_k) |i\rangle, \end{aligned} \tag{8}$$

$$\begin{aligned}\mathbf{J}^{\rm spin}_{\rm intr} &= \mathrm{i} \langle f| \sum_{i=1;\, j,k\neq i}^{3} \frac{1}{E_i + E'_i} \\ &\quad \times \Bigg[\frac{4}{9} e_i \boldsymbol{\sigma}_i \times (\mathbf{p}'_i - \mathbf{p}_i) \delta(\mathbf{p}'_i - \mathbf{p}_i - \mathbf{q}) \delta(\mathbf{p}'_j - \mathbf{p}_j) \delta(\mathbf{p}'_k - \mathbf{p}_k) \\ &\quad + \frac{1}{9} e_j \boldsymbol{\sigma}_j \times (\mathbf{p}'_j - \mathbf{p}_j) \delta(\mathbf{p}'_j - \mathbf{p}_j - \mathbf{q}) \delta(\mathbf{p}'_k - \mathbf{p}_k) \delta(\mathbf{p}'_i - \mathbf{p}_i) \\ &\quad + \frac{1}{9} e_k \boldsymbol{\sigma}_k \times (\mathbf{p}'_k - \mathbf{p}_k)\, \delta(\mathbf{p}'_k - \mathbf{p}_k - \mathbf{q}) \\ &\quad \times \delta(\mathbf{p}'_i - \mathbf{p}_i)\, \delta(\mathbf{p}'_j - \mathbf{p}_j) \Bigg] |i\rangle, \end{aligned} \tag{9}$$

where $\boldsymbol{\sigma}_i$ means that the matrix element is taken on the spin of the i-th quark.

The two-body part of the electromagnetic current operator can be derived directly from the continuity equation

$$\begin{aligned} &\mathbf{q} \cdot \Big(\mathbf{J}^{\rm cm}_{[1]} + \mathbf{J}^{\rm intr}_{[1]} + \mathbf{J}_{[2]} \Big) = \langle f| \Big[H, J^0_{[1]} \Big] |i\rangle = \\ &\frac{1}{E+E'} \langle f| \Big[\mathcal{M}_s^2 + R_s^2, J^0_{[1]} \Big] |i\rangle = \frac{1}{E+E'} \langle f| \Big[R_s^2, J^0_{[1]} \Big] |i\rangle \\ &+ \frac{2M}{E+E'} \Big(\langle f| \Big[\mathcal{M}_s - V, J^0_{[1]} \Big] + \Big[V, J^0_{[1]} \Big] |i\rangle \Big), \end{aligned} \tag{10}$$

consistently with the Fourier transform of the potential in eq. (1) and of the one-body charge operator. Here, we will consider only the $SU(2)$ sector of chiral symmetry, neglecting the strange quark. Therefore, the flavor (isospin) dependence of the charge generates non-vanishing exchange currents related to π and ρ exchanges only. In particular, the pseudoscalar piece gives the well-known isovector pion-pair ($\pi q\bar{q}$) and pion-in-flight ($\gamma\pi\pi$) currents [29]

Table 1. Parameters of the meson-exchange currents.

$m_{\rm u} = m_{\rm d}$	340 MeV	Λ_π	700 MeV
m_π	139 MeV	Λ_ρ	1200 MeV
m_ρ	770 MeV	$g_\pi^2/4\pi$	0.67
		$g_\rho^{\rm V\,2}/4\pi$	0.55
		$(g_\rho^{\rm V} + g_\rho^{\rm T})^2/4\pi$	1.31

$$\begin{aligned} &\mathbf{J}_{\pi q\bar{q}}(\mathbf{k}_i, \mathbf{k}_j) = \\ &i \frac{g_\pi^2}{4 m_i m_j} \left[\frac{\boldsymbol{\sigma}_i \cdot \mathbf{k}_i}{(\mathbf{k}_i^2 + m_\pi^2)} \boldsymbol{\sigma}_j \left(\frac{\Lambda_\pi^2 - m_\pi^2}{\mathbf{k}_i^2 + \Lambda_\pi^2} \right)^2 - (i \leftrightarrow j) \right] (\boldsymbol{\tau}_i \times \boldsymbol{\tau}_j)_z \end{aligned} \tag{11}$$

$$\begin{aligned} &\mathbf{J}_{\gamma\pi\pi}(\mathbf{k}_i, \mathbf{k}_j) = \\ &i \frac{g_\pi^2}{4 m_i m_j} \frac{\boldsymbol{\sigma}_i \cdot \mathbf{k}_i\, \boldsymbol{\sigma}_j \cdot \mathbf{k}_j}{(\mathbf{k}_i^2 + m_\pi^2)(\mathbf{k}_j^2 + m_\pi^2)} (\mathbf{k}_i - \mathbf{k}_j)\, (\boldsymbol{\tau}_i \times \boldsymbol{\tau}_j)_z \\ &\times \frac{(\Lambda_\pi^2 - m_\pi^2)^2}{(\mathbf{k}_i^2 + \Lambda_\pi^2)(\mathbf{k}_j^2 + \Lambda_\pi^2)} \left(1 + \frac{\mathbf{k}_i^2 + m_\pi^2}{\mathbf{k}_j^2 + \Lambda_\pi^2} + \frac{\mathbf{k}_j^2 + m_\pi^2}{\mathbf{k}_i^2 + \Lambda_\pi^2} \right), \end{aligned} \tag{12}$$

where $\mathbf{k}_i$, $\mathbf{k}_j$ are the momenta delivered to quarks i, j with mass $m_i = m_j = m_{\rm u} = m_{\rm d} = 340\,\mathrm{MeV}$ and the momentum conservation reads $\mathbf{q} = \mathbf{k}_i + \mathbf{k}_j$. The parameters $m_\pi, g_\pi, \Lambda_\pi$ are the mass, the coupling constant and the cut-off of the pion-quark vertex parametrized as

$$F(\mathbf{q}) = \frac{\Lambda^2 - m^2}{\Lambda^2 + \mathbf{q}^2}. \tag{13}$$

Analogously, the vector piece gives the well known isovector ρ-pair ($\rho q\bar{q}$) and ρ-in-flight ($\gamma\rho\rho$) currents [29]

$$\begin{aligned} &\mathbf{J}_{\rho q\bar{q}}(\mathbf{k}_i, \mathbf{k}_j) = i \frac{\big(g_\rho^{\rm V} + g_\rho^{\rm T}\big)^2}{4 m_i m_j} \\ &\left[\frac{\boldsymbol{\sigma}_i \times (\boldsymbol{\sigma}_j \times \mathbf{k}_j)}{(\mathbf{k}_j^2 + m_\rho^2)} \left(\frac{\Lambda_\rho^2 - m_\rho^2}{\mathbf{k}_j^2 + \Lambda_\rho^2} \right)^2 - (i \leftrightarrow j) \right] (\boldsymbol{\tau}_i \times \boldsymbol{\tau}_j)_z, \end{aligned} \tag{14}$$

$$\begin{aligned} &\mathbf{J}_{\gamma\rho\rho}(\mathbf{k}_i, \mathbf{k}_j) = \\ &\quad i \left[\big(g_\rho^{\rm V}\big)^2 + \frac{\big(g_\rho^{\rm V} + g_\rho^{\rm T}\big)^2}{4 m_i m_j} (\boldsymbol{\sigma}_i \times \mathbf{k}_i) \cdot (\boldsymbol{\sigma}_j \times \mathbf{k}_j) \right] \\ &\quad \times \frac{(\mathbf{k}_i - \mathbf{k}_j)}{(\mathbf{k}_i^2 + m_\rho^2)(\mathbf{k}_j^2 + m_\rho^2)} \frac{\big(\Lambda_\rho^2 - m_\rho^2\big)^2}{(\mathbf{k}_i^2 + \Lambda_\rho^2)(\mathbf{k}_j^2 + \Lambda_\rho^2)} \\ &\quad \times \left(1 + \frac{\mathbf{k}_i^2 + m_\rho^2}{\mathbf{k}_j^2 + \Lambda_\rho^2} + \frac{\mathbf{k}_j^2 + m_\rho^2}{\mathbf{k}_i^2 + \Lambda_\rho^2} \right) (\boldsymbol{\tau}_i \times \boldsymbol{\tau}_j)_z. \end{aligned} \tag{15}$$

The parameters $m_\rho, \Lambda_\rho, g_\rho^{\rm V}, g_\rho^{\rm T}$ are the ρ mass, cut-off, vector and tensor coupling constants of the ρ-quark vertex, respectively. All the parameter values have been kept the same as the ones used in refs. [2,3] for reproducing the baryon spectrum (see table 1).

Table 2. Contributions to the magnetic moments of the proton and neutron from different currents.

N		$\mu_{\mathrm{N}}^{[1]}$		$\mu_{\mathrm{N}}^{\pi q\bar{q}}$	$\mu_{\mathrm{N}}^{\gamma\pi\pi}$	$\mu_{\mathrm{N}}^{\rho q\bar{q}}$	$\mu_{\mathrm{N}}^{\gamma\rho\rho}$	$\mu_{\mathrm{N}}^{\rho\pi\gamma}$	μ_{N}	exp
p	u	1.773	1.988	−0.126	0.737	−0.109	0.202	0.048	2.740	2.793
	d	0.215								
n	u	−0.430	−1.316	0.126	−0.737	0.109	−0.202	0.048	−1.972	−1.913
	d	−0.886								

We have also explored the role of "model-dependent" two-body currents, namely of operators which are not constrained by the continuity eq. (10) because of their transverse nature. In particular, we have considered the well-known isoscalar $\rho\pi\gamma$ current

$$\mathbf{J}_{\rho\pi\gamma}(\mathbf{k}_i,\mathbf{k}_j) = \mathrm{i}\frac{g_\pi}{2m}\frac{g_\rho^{\mathrm{V}}}{m_\rho} g_{\rho\pi\gamma}(\mathbf{k}_i\times\mathbf{k}_j)(\boldsymbol{\tau}_i\cdot\boldsymbol{\tau}_j) \times\left[\frac{\boldsymbol{\sigma}_i\cdot\mathbf{k}_i}{\left(\mathbf{k}_i^2+m_\pi^2\right)\left(\mathbf{k}_j^2+m_\rho^2\right)}\frac{\Lambda_\rho^2-m_\rho^2}{\mathbf{k}_j^2+\Lambda_\rho^2}\frac{\Lambda_\pi^2-m_\pi^2}{\mathbf{k}_i^2+\Lambda_\pi^2}-(i\leftrightarrow j)\right], \tag{16}$$

where $g_{\rho\pi\gamma} = 0.578\pm0.028$ is deduced in accordance with the Vector Meson Dominance hypothesis (VMD) for the $\rho\to\pi\gamma$ decay width [30].

The magnetic moment is the global sum of the contributions corresponding to each previous component of the current operator:

$$\mu_{\mathrm{N}} = \mu_{\mathrm{N}}^{[1]} + \mu_{\mathrm{N}}^{\pi q\bar{q}} + \mu_{\mathrm{N}}^{\gamma\pi\pi} + \mu_{\mathrm{N}}^{\rho q\bar{q}} + \mu_{\mathrm{N}}^{\gamma\rho\rho} + \mu_{\mathrm{N}}^{\rho\pi\gamma}. \tag{17}$$

In table 2 the different results are shown. The one-body contribution is the leading one, as expected, but the proper treatment of the cm and intrinsic Hamiltonians represents a substantial improvement with respect to refs. [29,31]. The isovector two-body contributions, constrained by the continuity equation, show large cancellations but are globally important and act in opposite and correct ways according to the nucleon isospin. Finally, the isoscalar $\rho\pi\gamma$ contribution, though small, adds with the same sign both to proton and neutron magnetic moments, thus reducing the deviation of the theoretical values from the observed ones. The net theoretical result is in good agreement with the experiment, specifically with an error of about 1.5% for the proton and of about 2% for the neutron.

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