Flavor structure of the octet magnetic moments

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Abstract

We use the chiral quark-soliton model to identify all symmetry breaking terms linear in $m_s$ and investigate the strange magnetic moment in a “model-independent” way. Assuming hedgehog symmetry and employing the collective quantization, we obtain the most general expression for the flavor-singlet and flavor-octet magnetic moments in terms of seven independent parameters. Having fitted these parameters to the experimental magnetic moments of the octet baryons, we show that the strange magnetic moment turns out to be positive. The best fit obtained by minimizing $\chi^2$ assuming 15% theoretical accuracy yields: $\mu_s^{(s)} = (0.41 \pm 0.18) \mu_N$.

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The strangeness content of the nucleon has been a hot issue. The analysis of the \( \pi N \) sigma term \( \Sigma_{\pi N} \) implies a non negligible scalar density of the strange quark in the nucleon [1], once known as “\( \Sigma_{\pi N} \) puzzle”. In the axial channel, deep inelastic muon scattering conducted by European Muon Collaboration (EMC) almost ten years ago indicates that there is a sizeable contribution of the strange quark spin \( \Delta s \) to the proton spin, which was once called “spin crisis”. The measurement of elastic \( \nu p/\bar{\nu}p \) cross section at Brookhaven (BNL experiment 734), from which the strange axial form factor was extracted, came to the more or less same conclusion. In the vector channel the strange magnetic moment presents another key of understanding the strangeness content of the nucleon [3]. Great deal of theoretical effort has been put into the investigation of the strange magnetic moment and strange radius [4]–[17], which gave the range of possible values of the strange magnetic moment: \(-0.75 \rightarrow 2.2\) \( \mu_N \). This indicates that the strange properties of the nucleon in the vector channel are still poorly understood.

Very recently, the SAMPLE collaboration has announced the first measurement of the strange magnetic form factor \( G^s_M \) at \( Q^2 = 0.1 \text{ GeV}^2 \) [18]:

\[
G^s_M = (+0.23 \pm 0.37 \pm 0.15 \pm 0.19) \mu_N. \tag{1}
\]

Although the central value for \( G^s_M \) is positive, it is still too early to draw any firm conclusion, because of large experimental errors. One has to wait for other experiments which are presently either being conducted or planned at various electron accelerators [19–24]. It is therefore of importance to make theoretical predictions for this quantity.

Recently [25–27] magnetic moments of baryons have been studied within the framework of the chiral quark-soliton model (\( \chi QSM \)) – see review article [28] and references therein. In Ref. [27] the ”model-independent” way of analyzing magnetic moments has been proposed, in which the intrinsic dynamical parameters of the chiral quark-soliton model were fitted to the experimental data of the octet magnetic moments. In this way the magnetic moments of the baryon decuplet have been predicted. Actually, the results for the \( \mu_{\Omega^-} \) and \( \mu_{\Delta^{++}} \) are in a remarkable agreement with the data. Encouraged by these results, in the present work we apply the same procedure to predict the flavor magnetic moments. This requires the evaluation of the singlet part of the magnetic moment operator, which has not been calculated in Ref. [27].

Our present investigation is in line with the recent work by Hong, Park, and Min [17], which studied the strange magnetic moment of the nucleon within the framework of the chiral bag model using the ”model-independent” method analogous to that of Ref. [27]. However, in Ref. [17] the singlet magnetic moment does not contain the SU(3) symmetry breaking contribution; also the triplet and octet parts of the magnetic moments do not contain the symmetry breaking pieces which are present in our approach. Hence, we aim in this letter at investigating the flavor structure of the magnetic moments of the SU(3) baryons, taking consistently into account the contribution of the mass of the strange quark up to the linear order. We also discuss flavor structure of the magnetic moments for other members of the baryon octet.

Although it turns out that our numerical results for neutron and proton are very close to the ones of Ref. [17], we discuss in more detail the reliability of the whole approach by calculating the flavor magnetic moments for all baryons in the octet. We are confident that \( \mu^s_N \) is positive and of the order 0.2–0.6 \( \mu_N \).
The vector form factors of the SU(3) baryons are defined by the matrix elements of the U$_V$(3) vector currents:

$$
\langle B(p')|V^{(a)}_{\mu}|B(p)\rangle = \bar{u}_B(p') \left[ \gamma_\mu F_1^{(a)}(q^2) + i\sigma_{\mu\nu} \frac{q^\nu}{2M_N} F_2^{(a)}(q^2) \right] u_B(p) 
$$

with $a = 0, 3, 8$. The U$_V$(3) currents are defined as follows:

$$
V_\mu^{(0)} = \frac{1}{3} \gamma_\mu \gamma_5 \bar{\psi}, \quad V_\mu^{(3)} = \bar{\psi} \gamma_\mu \gamma_5 \gamma_5 \bar{\psi}, \quad V_\mu^{(8)} = \bar{\psi} \gamma_\mu \gamma_5 \psi.
$$

$q^2$ in Eq.(2) denotes the square of the four momentum transfer. The magnetic Sachs form-factor $G_M^{(a)}$ can be defined in terms the vector form-factors $F_1^{(a)}$ and $F_2^{(a)}$ in the following way:

$$
G_M^{(a)} = F_1^{(a)}(q^2) + F_2^{(a)}(q^2).
$$

In the non relativistic limit $G_M^{(a)}$ can be related to the time and space components of the U$_V$(3) vector currents, namely:

$$
\langle B(p')|V_i^{(a)}|B(p)\rangle = \frac{1}{2M_N} G_M^{(a)}(q^2)i\epsilon_{ijk}q^j \langle s'|\sigma_k|s\rangle,
$$

where $\sigma_k$ denotes Pauli spin matrices while $|s\rangle$ is the corresponding spin state of the baryon. The magnetic moments $\mu^{(a)}$ corresponding to the vector currents are identified with $G_M^{(a)}(0)$.

The matrix elements given above are related to the correlator

$$
\langle 0|J_B(x, T)\bar{\psi}\gamma_\mu \hat{Q}\psi J_B^\dagger(y, 0)|0\rangle
$$

at large Euclidean time $T$. The baryon current $J_B$ can be constructed from $N_c$ quark fields,

$$
J_B = \frac{1}{N_c!} \varepsilon^{i_1\ldots i_{N_c}} \Gamma^{\alpha_1\ldots \alpha_{N_c}}_{SS_3II_3Y} \psi_{\alpha_1 i_1} \ldots \psi_{\alpha_{N_c} i_{N_c}}
$$

$\alpha_1 \ldots \alpha_{N_c}$ are spin–isospin indices, $i_1 \ldots i_{N_c}$ are color indices, and the matrices $\Gamma^{\alpha_1\ldots \alpha_{N_c}}_{SS_3II_3Y}$ are taken to endow the corresponding current with the quantum numbers $SS_3II_3Y$. $J_B(J^\dagger_B)$ annihilates (creates) the baryon state at given time $T$. Taking into account the rotational corrections to the order $O(1/N_c)$ and $O(m_s)$ (for details, see for example Ref. [28]), the general expressions for the collective magnetic moment operators $\hat{\mu}^{(a)}$ can be written as follows:

$$
\hat{\mu}^{(0)} = \frac{1}{3} w_3 \hat{S}_3 + \frac{\sqrt{3}}{3} m_s (w_5 - w_6) D^{(8)}_{33},
$$

$$
\hat{\mu}^{(a)} = \left( w_1 + m_s w_1^2 \right) D^{(8)}_{33} + w_2 d_{pq3} D^{(8)}_{ap} \hat{S}_q + \frac{w_3}{\sqrt{3}} D^{(8)}_{as3} \hat{S}_3
$$

$$
+ m_s \left[ \frac{w_4}{\sqrt{3}} d_{pq3} D^{(8)}_{ap} D^{(8)}_{sq} + w_5 \left( D^{(8)}_{a3} D^{(8)}_{88} + D^{(8)}_{a8} D^{(8)}_{83} \right) + w_6 \left( D^{(8)}_{a3} D^{(8)}_{88} - D^{(8)}_{a8} D^{(8)}_{83} \right) \right]
$$

with $a = 3$ or 8. $\hat{S}_a$ stands for the operator of the generalized spin acting on the angular variable $R(t)$ [29]. $D^{(R)}_{ab}$ ($R$) denotes the SU(3) Wigner matrix in the representation $R$ and
\( m_s \) is the mass of the strange current quark. The parameters \( w_i \) depend on the dynamics of a specific hedgehog model.

Eq. (8) is different from Eq. (5) of Ref. [17] by terms proportional to \( m_s (w_5 - w_6) \). In Ref. [17] \( w_5 = w_6 \), and as a result the singlet magnetic moment operator \( \hat{\mu}^{(0)} \) is degenerated in the octet and decuplet. Also \( \hat{\mu}^{(3)} \) and \( \hat{\mu}^{(8)} \) get extra contributions in the present case, where \( w_5 \neq w_6 \). Numerically these terms are not very important in the case of the octet magnetic moments. This is, however, visible only \textit{a posteriori}. Magnetic moments are particularly sensitive to the SU(3) symmetry breaking, therefore careful analysis of all possible symmetry breaking terms is of quite importance.

The operator (8) has to be sandwiched between the octet collective wave functions. However, strictly speaking, octet baryon wave functions are no longer pure octet states:

\[
\Psi_B(R) = \Psi_B^{(8)}(R) + m_s c_{10}^{B} \Psi_B^{(10)}(R) + m_s c_{27}^{B} \Psi_B^{(27)}(R),
\]

where

\[
c_{10}^{B} = c_{10} \sqrt{5} \begin{bmatrix} 1 \\ 0 \\ 1 \\ 0 \end{bmatrix}, \quad c_{27}^{B} = c_{27} \begin{bmatrix} \sqrt{6} \\ 3 \\ 2 \\ \sqrt{6} \end{bmatrix}
\]

in the basis of \([N, \Lambda, \Sigma, \Xi]\). The coefficients \( c_{10} \) and \( c_{27} \) are defined in Ref. [27] and depend on the inertia parameters \( I_i \) and \( K_i \) \((i = 1, 2)\) calculable in the model. For the purpose of the present "model independent" analysis we shall use the following approximate equality:

\[
\frac{K_1}{I_1} \simeq \frac{K_2}{I_2}.
\]

Eq. (11) comes from the analysis of the mass splittings for the strange quark mass \( m_s = 180 \text{ MeV} \) [27]. It gets further support from the study of the model behavior in the limit of the small soliton mass, which corresponds to the non-relativistic quark model limit of the present model. In this limit equality (11) becomes exact.

Using (11) we can write \( c_{10} \) and \( c_{27} \) as follows:

\[
m_s c_{10} = c, \quad m_s c_{27} = \frac{3}{5} c,
\]

where \( c \) is a parameter corresponding to the moment of inertia \( I_2 \), which cannot be extracted from the mass splittings. Let us remind that \( I_2 \) enters into the splittings between octet and exotic baryonic states belonging to 10 or 27 [30].

The collective wave function can be explicitly written in terms of the SU(3) Wigner \( D^{(R)} \) function:

\[
\Psi_B^{(R)} = (-)^{s_3-1/2} \sqrt{\text{dim}(R)} [D_{(YTT)}^{(R)}(-1Ss_3)]^*.
\]

The flavor content of baryon magnetic moments is then expressed by
\[
\mu_B^{(u)} = \mu_B^{(0)} + \frac{1}{2} \mu_B^{(3)} + \frac{1}{2\sqrt{3}} \mu_B^{(8)}, \\
\mu_B^{(d)} = \mu_B^{(0)} - \frac{1}{2} \mu_B^{(3)} + \frac{1}{2\sqrt{3}} \mu_B^{(8)}, \\
\mu_B^{(s)} = \mu_B^{(0)} - \frac{1}{\sqrt{3}} \mu_B^{(8)}.
\]

Using Eqs. (8, 14), we can easily express the magnetic moments of the octet:

\[
\mu_B = \frac{1}{2} B \left( \mu^{(3)} + \frac{1}{\sqrt{3}} \mu^{(8)} \right) = \frac{1}{3} B \left( 2\mu^{(u)} - \mu^{(d)} - \mu^{(s)} \right).
\]

With the accuracy \(O(m_s)\), we can express flavor magnetic moments as follows:

\[
\mu_B^{(f)} = \mu_B^{(f)}(m_0^0) + \mu_B^{(f)}(m_s^0, \text{op}) + \mu_B^{(f)}(m_s^0, \text{wf}),
\]

where \(f\) equals \(u\), \(d\) or \(s\). \(\mu_B^{(f)}(m_0^0)\) denotes the SU(3) symmetric part of the magnetic moments, while the symmetry breaking parts \(\mu_B^{(f)}(m_s^0, \text{op})\) and \(\mu_B^{(f)}(m_s^0, \text{wf})\) correspond to the symmetry breaking in the operator and in the baryon wave functions, respectively.

As in Ref. [27], we introduce a new set of parameters:

\[
v = \frac{1}{60} (w_1 - \frac{1}{2} w_2), \quad w = \frac{1}{120} w_3, \\
\sqrt{x} = \frac{1}{540} m_s w_4, \quad y = \frac{1}{90} m_s w_5, \quad z = \frac{1}{30} m_s w_6, \\
p = \frac{1}{6} c \left( w_1 + w_2 + \frac{1}{2} w_3 \right), \quad q = -\frac{1}{150} c \left( w_1 + 2w_2 - \frac{3}{2} w_3 \right),
\]

where \(w_1 = w_1^1 + m_s w_1^2\). We can now express Eq. (14) explicitly in terms of the new variables of Eq. (17):

\[
\mu_p^{(u)} = \mu_p^{(d)} = -8v + 24w - 8x - 8y + 8q, \\
\mu_p^{(u)} = \mu_p^{(d)} = 6v + 22w + 14x + 2y + 2z + 2p + 4q, \\
\mu_{\Sigma^+}^{(u)} = \mu_{\Sigma^+}^{(d)} = 3v + 21w - 9x + 9y - 3z + 9q, \\
\mu_{\Sigma^+}^{(u)} = \mu_{\Sigma^+}^{(d)} = -8v + 24w - 4x - 10y + 4z + 4q, \\
\mu_{\Xi^0}^{(u)} = \mu_{\Xi^0}^{(d)} = -3v + 19w + 5x - 7y + 3z + p + 4q, \\
\mu_{\Xi^0}^{(u)} = \mu_{\Xi^0}^{(d)} = 2v + 14w + 14x - 4y + 2z + 2p + 4q, \\
\mu_{\Xi^-}^{(u)} = \mu_{\Xi^-}^{(d)} = 6v + 22w - 4x + 11y - 5z + 4q, \\
\mu_{\Xi^-}^{(u)} = \mu_{\Xi^-}^{(d)} = 2v + 14w - 8x - 7y - 3z + 8q, \\
\mu_{\Xi^0}^{(s)} = \mu_{\Xi^0}^{(d)} = 2v + 14w - 6x - 3y + z - 2p - 12q, \\
\mu_{\Lambda}^{(s)} = -6v + 18w + 18x + 9y - 3z - 18q, \\
\mu_{\Sigma^+}^{(s)} = \mu_{\Xi^0}^{(d)} = 6v + 22w - 10x - 13y + 3z - 2p - 8q, \\
\mu_{\Xi^-}^{(s)} = \mu_{\Xi^-}^{(d)} = -8v + 24w + 12x + 18y - 4z - 12q.
\]

(18)
Equations (18) exhibit both isospin and "hypercharge" symmetries of flavor magnetic moments.

In Ref. [27] we have fitted parameters $v$, $w$, $x$, $y$, $z$, $p$ and $q$ to the experimental data of the octet magnetic moments. It is of great convenience to fit first 2 parameters $v$ and $w$ in such a way that they are independent of $m_s$:

$$v = -0.268 = \frac{(2n + 3\Xi^0 + \Xi^- - p - 2\Sigma^- - 3\Sigma^+) }{60},$$
$$w = 0.060 = \frac{(3p + 4n + 3\Xi^0 - 3\Xi^- - 4\Sigma^- - \Sigma^+) }{60}. \quad (19)$$

This procedure has the advantage that one can study the chiral limit of the flavor magnetic moments, by tuning parameters $x, y, z, p$ and $q$ to zero. As already mentioned, throughout our analysis we assume that the mass of the strange quark is equal to 180 MeV.

Variable $z$ can also be extracted independently and the result reads:

$$z = -0.080 = \frac{1}{6}(n - \Sigma^- + \Sigma^+ - \Xi^0 - p + \Xi^-). \quad (20)$$

With these values for $v, w$ and $z$ we obtain 4 equations for the remaining 4 variables:

$$\begin{bmatrix} -8 & -5 & 0 & 8 \\ 14 & 5 & 2 & 4 \\ -9 & 0 & 0 & 9 \\ -4 & -1 & 0 & 4 \end{bmatrix} \begin{bmatrix} x \\ y \\ p \\ q \end{bmatrix} = \begin{bmatrix} 0.326 \\ -0.343 \\ 0.134 \\ 0.157 \end{bmatrix}, \quad \begin{bmatrix} p, \Xi^- \\ n, \Sigma^- \\ \Lambda \\ \Sigma^+, \Xi^0 \end{bmatrix}. \quad (21)$$

Four equations (21) are obtained by inserting (19) and (20) into the equations for the magnetic moments (18) and (15) (or equivalently into Eq.(50) of Ref. [27]). They are obtained from the equations for the magnetic moments of the particles given in the last column of Eq.(21)\(^2\). By redefining $x = x' - p/9$ and $q = q' - p/9$ one can get rid of the $p$ dependence.

So in fact we are left with 4 equations for 3 parameters: $x'$, $y$ and $q'$. In principle one could expect to obtain 4 independent fits by removing one equation from the set (21) and then solving it for $x'$, $y$ and $q'$. However, by removing the equation for $(n, \Sigma^-)$ we get a set of linearly-dependent equations, so we content ourselves with only three fits\(^3\) obtaining:

$$\begin{array}{ccc}
\text{fit(1)} & \text{fit(2)} & \text{fit(3)} \\
p, \Xi^- \text{ removed} & \Lambda \text{ removed} & \Sigma^+, \Xi^0 \text{ removed} \\
x' = 0.005, & x' = -0.026, & x' = -0.012, \\
y = -0.097, & y = -0.004, & y = -0.037, \\
q' = 0.020, & q' = 0.012, & q' = 0.003. \\
\end{array} \quad (22)$$

\(^1\)Strictly speaking $v$ and $w$ can be extracted independently of the remaining unknowns in Eq.(18), it is, however, impossible to get rid of a weak $m_s$ dependence in $w_1 = w_1^1 + m_s w_1^2$, which we neglect throughout this paper.

\(^2\)For example the first equation in (21) has been obtained from the equation for proton and it is identical to the one obtained from the equation for $\Xi^-$.

\(^3\)In Ref [27] we have presented results corresponding to fit (2).
The results, both for flavor and total magnetic moments are of course independent of $p$. They are displayed in Table I. The fits (1), (2) and (3) of Eq.(22) differ by $\chi^2$ which equals 2.78, 1.66 and 0.058 respectively. Fit (3) is by far the best; it misses $\Sigma^+$ by 2% and $\Xi^0$ by 7%. Guided by this accuracy we have minimized $\chi^2$ assuming that the average error is 15%. One should interpret this error as the theoretical error corresponding to the truncation of the perturbation series in $m_s$. The result reads:

$$ v = -0.266 \pm 0.030, \quad w = 0.064 \pm 0.036, \quad x' = -0.008 \pm 0.017, $$

$$ y = -0.050 \pm 0.042, \quad z = 0.093 \pm 0.190, \quad q' = 0.006 \pm 0.020, $$

(23)

corresponding to $\chi^2 = 0.017$. Errors have been calculated by allowing $\chi^2$ to change by 1. Results of fit (4) are given in Table I. In Table I we also present the model calculation ($\chi$QSM) for the constituent mass $M = 420$ MeV and $m_s = 180$ MeV and (where available) the results of the model independent fit 1 of Ref. [17] (HPM).

It should be mentioned that the non-relativistic limit of the chiral quark–soliton model (which, to some extent, can be used as a guiding line) predicts $w/v = -1/7, 3y/v = z/v = 8/21 m_s I_2$, and $q' = x' = 0$ which is in a qualitative agreement with our fit eq. (23). With values of the constants in non-relativistic limit we get:

$$ \frac{\mu_p}{\mu_n} = -\frac{3}{2} \left(1 - \frac{8}{135} m_s I_2\right), \quad \mu_s = 0, $$

(24)

these are the well known results of the $SU(6)$ non-relativistic quark model with particular mass correction to them (note that mass corrections for $\mu_s$ are exactly zero in the non-relativistic limit).

Before we comment on the results presented in Table I let us discuss the sensitivity of our fits to the symmetry breaking. To do this we reconsider fits (1), (2) and (3) and restore the linear $m_s$ dependence of the variables $x', y, z, and q'$. At this point we can profit from our fitting procedure in which chiral limit parameters $v$ and $w$ were fitted independently of $m_s$. The results for fit (3) are plotted in Fig.1. We can see from Fig.1 that $\mu_B^{(s)}$ is rather stable, as far as $m_s$ dependence is concerned. “Error bars” in Fig.1. correspond to the range of values for $\mu_B^{(s)} (m_s = 180$ MeV) for all fits (with fit (4) included). However, as can be seen from Fig.1, the nucleon line is almost flat: $\mu_N^{(s)}$ is quite insensitive to the symmetry breaking.

Of course the value $\mu_N^{(s)} = 0.41 \mu_B$ of fit (4) has its own error, which can be estimated by varying parameters $v, w, x', y, z$ and $q'$ within the error bars given in Eq.(23). The average error obtained by this procedure equals to $\pm 0.18$ assuming 15% accuracy in the $\chi^2$ fit (or $\pm 0.12$ for 10% accuracy).

Despite the uncertainties for the strange particles we conclude that the data suggest rather large strangeness contribution to the magnetic moment of the nucleon. On the contrary, the self-consistent $\chi$QSM gives for $\mu_p^{(s)}$ a small value, which is close to zero, but positive $^4$. This is due to the fact that parameter $w_3$ (or equivalently $w$) is much smaller

$^4$Note that Ref. [15] includes a numerical error in the calculation of the strange magnetic moment.
than the one obtained by our fitting procedure. As a result the “starting point” values for magnetic moments (at \( m_s = 0 \)) differ from the ones obtained by all our “model independent” fits with the help of Eq.(19). In the self-consistent \( \chi \)QSM linear \( m_s \) dependence is not strong enough to change the chiral limit results enough to reproduce the experimental data for the magnetic moments. Also the strange magnetic moment of the nucleon stays close to the SU(3) symmetry value which is almost zero (but negative).

Our present “model-independent” results for \( \mu_{p,n}^{(u)} \) and \( \mu_n^{(s)} \) agree quite well with the ones of Ref. [17]. Numerically the contribution of the terms proportional to the difference \( w_5 - w_6 \) is quite small; in the worse case of fit (1) it does not exceed 8 % of the proton magnetic moment.

In summary, we have investigated the flavor structure of the octet magnetic moments within the framework of the chiral quark-soliton model, employing a “model-independent” approach. The strange magnetic moment of the nucleon turns out to be positive. The effect of the SU(3) symmetry breaking is almost negligible in the case of the nucleon. The “model-independent” fit \( (\mu_N^{(s)} \approx (0.25 - 0.37)\, \mu_N) \) is stable and yields quite larger value than that of the self-consistent model calculation \( (\mu_N^{(s)} = 0.03\mu_N) \). On the other hand the best fit obtained by minimizing \( \chi^2 \) assuming 15% theoretical accuracy yields \( \mu_N^{(s)} = (0.41 \pm 0.18)\, \mu_N \).

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REFERENCES

Table 1: Magnetic moments and their flavor decomposition for 4 fits described in the text, for model calculations ($\chi_{QSM}$) and for fit 1 of Ref. [17](HPM).

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<th>$\mu^{(s)}$</th>
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