Magnetic Moments of the SU(3) Octet Baryons in the 
semibosonized SU(3) Nambu-Jona-Lasinio Model

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(August, 1995)

Abstract

We investigate the magnetic moments of the SU(3) octet baryons in the 
framework of the SU(3) semibosonized Nambu–Jona–Lasinio model. The ro-
tational 1/\(N_c\) corrections and strange quark mass in linear order are taken 
to account. We derive general relations between magnetic moments of the 
SU(3) octet baryons, based on the symmetry of our model. These relations 
indicate that higher order corrections such as \(O(m_s/N_c)\) and \(O(m_s^2)\) are 
relatively small. The magnetic moments of the octet baryons predicted by our 
model are quantitatively in a good agreement with experimental results within

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about 15\%.
I. INTRODUCTION

The semibosonized Nambu-Jona-Lasinio model (NJL) (often is called as the chiral quark soliton model) [1] is very successful in describing the static properties of the nucleon such as the mass splitting of the nucleon and Δ isobar, axial constants [2], magnetic moments [3], electromagnetic form factors [3], and so on. Recently, Blotz et al. [4] and Weigel et al. [5] showed that the SU(3) version of the model explains the correct mass splitting of the SU(3) octet and decuplet baryons. The model could also reproduce the axial constants \( g_A^{(0)} \), \( g_A^{(3)} \) and \( g_A^{(8)} \) with a remarkable agreement with experiments. In particular, the finding of the non-commutivity of the collective operators arising from the time-ordering enabled the model to solve the long-standing problem of the underestimate of the axial coupling constants and nucleon magnetic moments [2] in hedgehog models.

In the semibosonized NJL model, the baryon can be understood as \( N_c \) valence quarks coupled to the polarized Dirac sea bound by a nontrivial chiral background field in the Hartree approximation. The proper quantum numbers of baryons are obtained by the semi-classical quantization performed by integrating over zero-mode fluctuations of the pion field around the saddle point. The merit of the model is to interpolate between the naive quark model and the Skyrme model, which enables us to study the interplay between these two models.

In the present work, we shall investigate the magnetic moments of the SU(3) octet baryons in the SU(3) NJL model. Since the magnetic moments of the SU(3) octet baryons are experimentally well known, it is a good check for the model to reproduce them. Furthermore, we shall show that the model reaches the upper limit of the accuracy which can be attained in any model with “hedgehog symmetry”.

The outline of the present work is as follows: In the next section, we briefly describe the semibosonized SU(3) NJL model and show how to obtain the magnetic moments in the model. In section 3, we derive the general relations between magnetic moments of the octet baryons using the symmetry of the model and confront them with experimental data.
We show that subleading $O(m_s/N_c)$ and $O(m_s^2)$ corrections are relatively small, whereas $O(1/N_c)$ ones are fairly large. In section 4, we discuss the numerical results. We summarize the present work and draw conclusion in section 5.

II. FORMALISM

The semibosonized NJL model is described by a partition function in Euclidean space given by the functional integral over pseudoscalar meson and quark fields:

$$Z = \int \mathcal{D}\Psi \mathcal{D}\Psi^\dagger \mathcal{D}\pi^a \exp \left( - \int d^4x \Psi^\dagger iD\Psi \right),$$

$$= \int \mathcal{D}\pi^a \exp (-S_{\text{eff}}[\pi]),$$

where $S_{\text{eff}}$ is the effective action

$$S_{\text{eff}}[\pi] = -\text{Sp} \log iD.$$  (2)

$iD$ represents the Dirac differential operator

$$iD = \beta(-i\partial + \hat{m} + MU).$$  (3)

with the pseudoscalar chiral field

$$U = \exp (i\pi^a \lambda^a \gamma_5).$$  (4)

$\hat{m}$ is the matrix of the current quark mass given by

$$\hat{m} = \text{diag}(m_u, m_d, m_s) = m_0 1 + m_8 \lambda_8.$$  (5)

$\lambda^a$ represent the usual Gell-Mann matrices normalized as $\text{tr}(\lambda^a \lambda^b) = 2\delta^{ab}$. Here, we have assumed the isospin symmetry. $M$ stands for the dynamical quark mass arising from the spontaneous chiral symmetry breaking, which is in general momentum-dependent [6]. We regard $M$ as a constant and employ the proper-time regularization for convenience. The $m_0$ and $m_8$ in eq. (5) are respectively defined by
\[ m_0 = \frac{m_u + m_d + m_s}{3}, \quad m_8 = \frac{m_u + m_d - 2m_s}{2\sqrt{3}}. \] (6)

The operator \( iD \) is expressed in Euclidean space in terms of the Euclidean time derivative \( \partial_\tau \) and the Dirac one–particle hamiltonian \( H(U) \)

\[ iD = \partial_\tau + H(U) + \beta \hat{m} - \beta \bar{m} \mathbf{1} \] (7)

with

\[ H(U) = \frac{\bar{\alpha} \cdot \nabla}{\imath} + \beta MU + \bar{m} \mathbf{1}. \] (8)

\( \bar{m} \) is defined by \( (m_u + m_d)/2 = m_u = m_d \). \( \beta \) and \( \bar{\alpha} \) are the well–known Dirac hermitian matrices. Note that the NJL model is a low-energy effective model of QCD. Hence, the effective action given by eq.(1) can include higher order mass terms like \( \hat{m}^2 \Psi^\dagger \Psi \). However, the coefficient in front of \( \hat{m}^2 \Psi^\dagger \Psi \) is not theoretically known\(^1\). To go beyond the linear order of mass corrections, one should justify such a higher order term. Otherwise, it is meaningless to consider higher order mass corrections in the expansion of the quark mass. Therefore, we shall take into account the mass corrections only up to the linear order.

Many physical processes (semileptonic decays, electromagnetic transitions, electromagnetic form factors, etc.) are described by the one-current baryon matrix element:

\[ \langle B_2 | \bar{\psi} \Gamma \hat{O} \psi(x) | B_1 \rangle, \] (9)

where \( \Gamma = (\gamma_\mu, \gamma_\mu \gamma_5, \sigma_{\mu\nu}, \gamma_5) \) is a particular Dirac matrix depending on the physical observable and \( O \) is a \( SU(3) \) flavor matrix. For example, the matrix element in eq.(10) with

\[ \Gamma = \gamma_\mu, \quad \hat{O} = \frac{1}{2} \lambda^3 + \frac{1}{2\sqrt{3}} \lambda^8 \] (10)

is relevant to the electromagnetic form factors of the octet baryons (magnetic moments, electromagnetic square radii, etc.). This particular matrix element is a subject of the present paper.

\(^1\)The coefficient \( \hat{m} \Psi^\dagger \Psi \) is determined by the soft pion theorem.
One can relate the hadronic matrix element eq. (9) to a correlation function:

\[ \langle 0 | J_{B_1} (\vec{x}, T) \bar{\psi} \Gamma O \psi J_{B_2} (\vec{y}, 0) | 0 \rangle \] (11)

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

\[ J_B = \frac{1}{N_c} \epsilon^{i_1 \ldots i_{N_c}} \Gamma_{J_3 I_3 Y} \psi_{\alpha_1 i_1 \ldots \psi_{\alpha_{N_c} i_{N_c}}} \] (12)

\( \alpha_1 \ldots \alpha_{N_c} \) are spin–isospin indices, \( i_1 \ldots i_{N_c} \) are color indices, and the matrices \( \Gamma_{J_3 I_3 Y}^{\alpha_1 \ldots \alpha_{N_c}} \) are chosen in such a way that the quantum numbers of the corresponding current are equal to \( J J_3 I_3 Y \). \( J_B (J_B^\dagger) \) annihilates (creates) a baryon at large \( T \). The general expression for the matrix elements eq. (9) was derived in Ref. [23] with linear \( m_s \) corrections taken into account:

\[ \langle B_2 | \bar{\psi} \Gamma O \psi (x) | B_1 \rangle = -N_c \int d^3 x e^{iqx} \int \frac{d\omega}{2\pi} \text{tr} \left( x \frac{1}{\omega + iH} \gamma_4 \Gamma \lambda^A \right) \]
\[ \times \int dR \Psi_{B_2}(R) \Psi_{B_1}(R) \frac{1}{2} \text{tr} \left( R^\dagger \lambda^A \Gamma \lambda^B (R) \right) \]
\[ + iN_c \int d^3 x e^{iqx} \int \frac{d\omega}{2\pi} \text{tr} \left( x \frac{1}{\omega + iH} \gamma_4 \Gamma \lambda^A \right) \]
\[ \times \int dR \Psi_{B_2}(R) \Psi_{B_1}(R) \frac{1}{2} \text{tr} \left( R^\dagger \lambda^A \Gamma \lambda^B (R) \right), \] (13)

where \( q \ll M_N \) is the momentum transfer and \( \lambda^A = \left( \sqrt{\frac{2}{3}} \mathbf{1}, \lambda^a \right) \). In eq. (13) a regularization is not shown for simplicity (see Ref. [24] for details). \( \Psi_B (R) \) are the rotational wave functions of the baryon. \( \Psi_B (R) \) requires the corrections due to the strange quark mass \( (m_s) \), since we treat the \( m_s \) perturbatively. Hence, \( \Psi_B (R) \) can be written by

\[ \Psi_B (R) = \Psi_B^{(8)} (R) + c_B^{10} \Psi_B^{(10)} (R) + c_B^{27} \Psi_B^{(27)} (R) \] (14)

with

\[ c_B^{10} = \frac{\sqrt{5}}{15} (\sigma - r_1) \begin{bmatrix} 1 \\ 0 \\ 1 \\ 0 \end{bmatrix} I_2 m_s, \quad c_B^{27} = \frac{1}{75} (3\sigma + r_1 - 4r_2) \begin{bmatrix} \sqrt{6} \\ 3 \\ 2 \\ \sqrt{6} \end{bmatrix} I_2 m_s. \] (15)

Here, \( B \) denotes the SU(3) octet baryons with the spin 1/2. The constant \( \sigma \) is related to the \( \pi N \) sigma term \( \Sigma = 3/2 (m_u + m_d) \sigma \) and \( r_i \) designates \( K_i / I_i \), where \( K_i \) stands
for the anomalous moments of inertia defined in Ref. [4]. Recently, [2,8,9] the rotational $1/N_c$ corrections for matrix elements of vector and axial currents were derived, the general expression (without any regularization) for these corrections has a form:

$$
\Delta \Omega^i \langle B_2 | \bar{\psi} \Gamma O \psi (x) | B_1 \rangle = iN_c \int d^3 x e^{iqx} \int \frac{d\omega}{2\pi} \int \frac{d\omega'}{2\pi} \frac{1}{P} \frac{1}{\omega - \omega'} (I^{-1})_{aa'}
\times \text{tr} \langle x | \frac{1}{\omega + iH} \lambda^{a'} \frac{1}{\omega' + iH} \gamma_4 \Gamma \lambda^b | x \rangle
\times f^{abc} \int dR \bar{\Psi}_{B_2} (R) \Psi_{B_1} (R) \frac{1}{2} \text{tr} (R^i \lambda^c RO)
+ N_c \int d^3 x e^{iqx} \int \frac{d\omega}{2\pi} \text{tr} \langle x | \frac{1}{\omega + iH} \lambda^a \frac{1}{\omega' + iH} \gamma_4 \Gamma \lambda^b | x \rangle
\times \int dR \bar{\Psi}_{B_2} (R) \frac{1}{2} \{ \text{tr} (R^i \lambda^c RO), \hat{\Omega}^a \} + \Psi_{B_1} (R). \quad (16)
$$

Where $I_{ab}$ is a matrix of the moments of inertia for the soliton, $\hat{\Omega}^a$ is an operator of angular velocities acting on angular variables $R$ (details can be found in [4]). In what follows we shall use these expressions to calculate the magnetic moments of the SU(3) octet baryons.

### III. MAGNETIC MOMENTS IN THE MODEL

Using the general expressions eq.(13) and eq.(16) for the one current baryonic matrix elements, we can express the magnetic moments of the SU(3) octet baryons with the $m_s$ and rotational $1/N_c$ corrections in terms of the dynamic quantities $v_i$ depending on the concrete dynamics of the chiral quark soliton:

$$
\mu_B = v_1 \langle B | D_{Q3}^{(8)} | B \rangle + \frac{v_2}{N_c} d_{ab3} \langle B | D_{Qa}^{(8)} \cdot \hat{J}_b | B \rangle
+ m_s \left[ (v_3 d_{ab3} + v_4 S_{ab3} + v_5 F_{ab3}) \cdot \langle B | D_{Q4}^{(8)} D_{sb}^{(8)} | B \rangle \right]. \quad (17)
$$

Here we have introduced $SU(2)_T \times U(1)_Y$ invariant tensors

$$
d_{abc} = \frac{1}{4} \text{tr} (\lambda_a \{ \lambda_b, \lambda_c \} +),
S_{ab3} = \frac{1}{\sqrt{3}} (\delta_{a3} \delta_{b8} + \delta_{b3} \delta_{a8}),
$$

and

$$
F_{ab3} = \frac{1}{\sqrt{3}} (\delta_{a3} \delta_{b8} - \delta_{b3} \delta_{a8}). \quad (18)
$$
\[ Q = \frac{1}{2} \lambda^3 + \frac{1}{2\sqrt{3}} \lambda^8 \] stands for the charge operator in SU(3) flavor space. The rotational wave functions \( |B\rangle \) are given by eq. (14). The dynamic quantities \( v_i \) are independent of the hadrons involved. They have a general structure like:

\[
\sum_{m,n} \langle n|O_1|m\rangle \langle m|O_2|n\rangle f(E_n, E_m, \Lambda),
\]

where \( O_i \) are spin-isospin operators changing the grand spin of states \( |n\rangle \) by 0 or 1 and the double sum runs over all the eigenstates of the quark hamiltonian in the soliton field. The numerical technique for calculating such a double sum has been developed in [4,22,25].

Before we calculate the magnetic moments numerically, let us estimate the importance of \( 1/N_c \) corrections and the relative size of subleading \( O(m_s/N_c) \) corrections. To this end we employ a dynamically independent relations between magnetic moments arising from the “hedgehog” symmetry of the model. Hyperon magnetic moments are parametrized (in our approximation) by six parameters \( (v_1, v_2, v_3, v_4, v_5 \text{ and one parameter is contained in the rotational wave functions}) \). Looking upon them as free parameters, we obtain the relations between the hyperon magnetic moments and the magnetic moment of the \( \Sigma^0 \Lambda \) transition:

\[
\mu_{\Sigma^0} = \frac{1}{2}(\mu_{\Sigma^+} + \mu_{\Sigma^-}),
\]

\[
\mu_{\Lambda} = \frac{1}{12}(-12\mu_p - 7\mu_n + 7\mu_{\Sigma^-} + 22\mu_{\Sigma^+} + 3\mu_{\Xi^-} + 23\mu_{\Xi^0})
\times (1 + O(\frac{m_s}{N_c}) + O(m_s^2))
\]

\[
\mu_{\Sigma^0\Lambda} = -\frac{1}{\sqrt{3}}(-\mu_n + \frac{1}{4}(\mu_{\Sigma^+} + \mu_{\Sigma^-}) - \mu_{\Xi^0} + \frac{3}{2}\mu_{\Lambda} \cdot (1 + O(\frac{m_s}{N_c}) + O(m_s^2)),
\]

and one additional relation if we neglect rotational \( 1/N_c \) corrections, i.e. put \( v_2 = 0 \) in eq.(17):

\[
\mu_{\Xi^0} = (-3\mu_p - 4\mu_n + 4\mu_{\Sigma^-} + \mu_{\Sigma^+} + 3\mu_{\Xi^-}) \cdot (1 + O(\frac{1}{N_c}) + O(m_s^2)).
\]

Note that the analogous relations between hyperon magnetic moments was obtained by Adkins and Nappi [26] but they did not take into account mass corrections to the rotational baryon wave functions and neglected \( 1/N_c \) corrections. The first relation eq.(20) is trivially
fulfilled. It is an isospin relation and so has no corrections in both $1/N_c$ and $m_s$. The next two relations eq.(21) and eq.(22) empirically gives:

$$-(0.613 \pm 0.004) = -(0.402 \pm 0.10)$$

and

$$-(1.61 \pm 0.08) = -(1.48 \pm 0.03)$$

respectively, whereas the fourth relation eq.(23) gives:

$$-(1.250 \pm 0.015) = -(4.8 \pm 0.2).$$

We see that the fourth relation eq.(23) in which we neglect $1/N_c$ corrections is badly reproduced by experiment whereas the relations given by eqs.(20,21,22) seem to be successful. The explanation of this difference lies in different large $N_c$ properties of the relations. These relations have, in principle, corrections of order $O(1/N_c)$, $O(m_s/N_c)$ and $O(m_s^2)$, but in (20,21,22) all corrections proportional to any power of $1/N_c$ are cancelled. Hence the relations eqs. (20,21,22) are satisfied with the accuracy of the order $O(m_s/N_c)$, while the eq.(23) is gratified with the accuracy of $O(1/N_c)$. From these estimates one can conclude that the corrections of order $O(1/N_c)$ to magnetic moments numerically are large whereas those of the order $O(m_s/N_c)$ can be relatively small. These estimates provide us a lower limit for the systematic errors of computations in any “hedgehog” model for baryons by neglecting the non-computed $O(m_s/N_c)$ and $O(m_s^2)$ corrections, since any “hedgehog” model fulfills eqs. (20,21,22) which are deviated from the experiment by about 15%. Hence such a kind of models can not reproduce the experimental data of magnetic moments better than the above–mentioned limit of 15%. We shall see that in the NJL soliton model the accuracy for the hyperon magnetic moments is very close to its upper limit.

IV. NUMERICAL RESULTS AND DISCUSSION

In order to calculate eq.(17) numerically, we follow the well-known Kahana and Ripka method [12].
In table 1 we show the dependence of the magnetic moments of the SU(3) octet baryons on the constituent quark mass in the chiral limit \((m_s = 0)\). Both of the leading term and the rotational \(1/N_c\) corrections tend to decrease as the constituent quark mass \(M\) increases. In this limit the \(U\)-spin symmetry is not broken, so that we have the relations

\[
\mu_p = \mu_{\Sigma^+}, \quad \mu_n = \mu_{\Xi^0}, \\
\mu_{\Sigma^-} = \mu_{\Xi^-}, \quad \mu_{\Sigma^0} = -\mu_{\Lambda}.
\]  

Compared to the SU(2) results, the prediction of the SU(3) model \((m_s = 0)\) for the nucleon is different and seems to be better. It is due to the fact that in our approach the nucleon possesses the polarized hidden strangeness [29,30].

The rotational \(1/N_c\) corrections are equally important to the other octet members as shown in Table 1. As a result, the total rotational \(1/N_c\) corrections contribute to the magnetic moments around 50%.

The symmetry breaking terms, proportional to \(m_s\), lift the \(U\)-spin symmetry. The \(m_s\) corrections arise from the explicit dependence of the baryon matrix elements on the strange quark mass \(m_s\) (second term of eq.(13)) on the one hand, and on the other hand come from the solitonic rotational wave functions (details see in Refs. [29,8]). The latter correction appears in each column of Table 2 and is equally important as the former one.

It is interesting to compare the NJL model with the Skyrme model, since these two models are closely related. As Ref. [8] already made a comparison between the NJL model and the Skyrme model in case of the \(g_A\). Apparently both models have the same collective operator structures (see eq.(17)), but the origin of parameters \(v_i\) given in eq.(17) is quite distinct each other. In the NJL model, the coefficients \(v_i\) include the contribution from the noncommutivity of the collective operators [2] while it is absent in the Skyrme model, since the lagrangian of the Skyrme model is local in contrast to that of the semibosonized NJL. The coefficient \(v_2\) in the Skyrme model comes from the pseudoscalar mesons dominated by the induced kaon fluctuations. It is interesting to note that the Skyrme model needs explicit vector mesons in addition to pseudoscalar ones [10] in order to achieve the same algebraic
structure of the collective hamiltonian as it is obtained in the semibosonized NJL model with pseudoscalar mesons alone. Due to the introduction of vector mesons, it is inevitable to import large numbers of parameters into the Skyrme model. In table 2, we compare our results with the Skyrme model [10].

In Fig. 1 we show how large the predicted magnetic moments deviate from the experimental data. We observe that the $U$-spin symmetry is lifted almost equidistantly. It is due to the fact that we have only taken into account the linear order of the $m_s$ corrections. However, it is not theoretically justified to consider higher $m_s$ corrections as discussed briefly in section 2.

On the whole, the magnetic moments are in a good agreement with the experimental data within about 15%.

V. SUMMARY AND CONCLUSION

We have studied the magnetic moments of the SU(3) octet baryons in the framework of the semibosonized SU(3) NJL model, taking into account the rotational $1/N_c$ corrections and linear $m_s$ corrections. The only parameter we have in the model is the constituent quark mass $M$ which is fixed to $M = 420$ MeV by the mass splitting of the SU(3) baryons. We have shown that the NJL model reproduces the magnetic moments of the SU(3) octet baryons within about 15 %. The accuracy we have reached is more or less the upper limit which can be attained in any model with “hedgehog symmetry”.

ACKNOWLEDGEMENT

We would like to thank Christo Christov, Michal Praszalowicz and T. Watabe for helpful discussions. This work has partly been supported by the BMFT, the DFG and the COSY-Project (Jülich).
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TABLES

TABLE I. The dependence of the magnetic moments of the SU(3) octet baryons on the constituent quark mass $M$ without $m_s$ corrections: $\mu(\Omega^0)$ corresponds to the leading order in the rotational frequency while $\mu(\Omega^1)$ includes the subleading order.

| Baryon | 370 MeV | 420 MeV | 450 MeV | Exp |
|--------|---------|---------|---------|-----|
|        | $\mu_B(\Omega^0)$ | $\mu_B(\Omega^1)$ | $\mu_B(\Omega^0)$ | $\mu_B(\Omega^1)$ | $\mu_B(\Omega^0)$ | $\mu_B(\Omega^1)$ |     |
| $p$    | 1.13    | 2.57    | 1.01    | 2.27 | 0.90 | 2.11 | 2.79 |
| $n$    | -0.84   | -1.77   | -0.75   | -1.55 | -0.67 | -1.44 | -1.91 |
| $\Lambda$ | -0.42   | -0.88   | -0.38   | -0.78 | -0.34 | -0.72 | -0.61 |
| $\Sigma^+$ | 1.12    | 2.57    | 1.01    | 2.27 | 0.90 | 2.11 | 2.46 |
| $\Sigma^0$ | 0.42    | 0.88    | 0.38    | 0.78 | 0.34 | 0.72 |     |
| $\Sigma^-$ | -0.28   | -0.81   | -0.25   | -0.71 | -0.22 | -0.67 | -1.16 |
| $\Xi^0$ | -0.84   | -1.77   | -0.75   | -1.55 | -0.67 | -1.44 | -1.25 |
| $\Xi^-$ | -0.28   | -0.81   | -0.25   | -0.71 | -0.22 | -0.67 | -0.65 |
The magnetic moments of the SU(3) octet baryons predicted by our model are compared with the evaluation from the Skyrme model of Park and Weigel [10]. The experimental values are taken from Ref. [27]. The constituent quark mass is fixed as $M = 420$ MeV. The $\mu_B(\Omega^1, m_s)$ include subleading orders in $\Omega$ and $m_s$, which are our final values.

| Baryons | $\mu_B(\Omega^0, m_s^0)$ | $\mu_B(\Omega^1, m_s^0)$ | $\mu_B(\Omega^1, m_s^1)$ | Park & Weigel | Exp. |
|---------|------------------------|------------------------|------------------------|--------------|------|
| $p$     | 1.03                   | 2.29                   | 2.39                   | 2.36         | 2.79 |
| $n$     | -0.90                  | -1.69                  | -1.76                  | -1.87        | -1.91|
| $\Lambda$ | -0.35              | -0.75                  | -0.77                  | -0.60        | -0.61|
| $\Sigma^+$ | 1.02               | 2.28                   | 2.42                   | 2.41         | 2.46 |
| $\Sigma^0$ | 0.31               | 0.72                   | 0.75                   | 0.66         | —    |
| $\Sigma^-$ | -0.40              | -0.85                  | -0.92                  | -1.10        | -1.16|
| $\Xi^0$  | -0.74                  | -1.54                  | -1.64                  | -1.96        | -1.25|
| $\Xi^-$  | -0.23                  | -0.69                  | -0.68                  | -0.84        | -0.65|
| $|\Sigma^0 \rightarrow \Lambda|$ | 0.74               | 1.42                   | 1.51                   | 1.74         | 1.61 |
Figure captions

**Fig. 1:** The magnetic moments of the SU(3) octet baryons predicted by the semibosonized NJL model. The first column denoted by (1) shows the magnetic moments of the SU(3) octet baryons in case of $m_s = 0$ MeV. Due to the $U$-spin symmetry, those of corresponding baryons are degenerated. The second column denoted by (2) designates the case of $m_s = 180$ MeV. The dotted lines show the breaking of the $U$-spin symmetry due to large $m_s$. The third column by (3) is for the experimental data. The constituent quark mass $M = 420$ MeV is chosen for our theoretical results.
(1) \( m_s = 0 \)  \hspace{1cm} (2) \( m_s = 180 \)  \hspace{1cm} (3) \( \text{Exp} \)