Electromagnetic form factors of octet baryons with the nonlocal chiral effective theory

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The electromagnetic form factors of octet baryons are investigated with the nonlocal chiral effective theory. The nonlocal interaction generates both the regulator which makes the loop integral convergent and the $Q^2$ dependence of form factors at tree level. Both octet and decuplet intermediate states are included in the one loop calculation. The momentum dependence of baryon form factors are studied up to 1 GeV$^2$ with the same number of parameters as for the nucleon form factors. The obtained magnetic moments of all the baryon octets as well as the radii are in good agreement with the experimental data and/or lattice simulation.
I. INTRODUCTION

The study of electromagnetic form factors of hadrons is of crucial importance to understand their sub-structure. A lot of theoretical and experimental efforts have been made in this field. On the one hand, with the upgrade of the experimental facility, the parton distribution functions (PDFs) from the deep inelastic scattering as well as the form factors at relatively large momentum transfer from the elastic scattering can be extracted [1, 2]. On the other hand, many measurements on form factors have been carried out at very small momentum transfer to get the information of nucleon radii as accurate as possible [3, 4].

Theoretically, though QCD is the fundamental theory to describe strong interactions, it is difficult to study hadron physics using QCD directly. There are many phenomenological models, such as the cloudy bag model [5], the constituent quark model [6], the 1/Nc expansion approach [7], the Nambu-Jona-Lasinio (NJL) model [8], the perturbative chiral quark model (PCQM) [9], the extended vector meson dominance model [10], the SU(3) chiral quark model [11], the quark-diquark model [12], etc. These model calculations are helpful to provide the physical scenario for the hadron structure.

Besides the phenomenological quark models, there are two systematic methods in hadron physics. One is the lattice simulation and the other is effective field theory (EFT) or chiral perturbation theory (ChPT). Historically, most formulations of ChPT are based on dimensional or infrared regularization. Though ChPT is a successful approach, for the nucleon electromagnetic form factors, it is only valid for $Q^2 < 0.1$ GeV$^2$ [13]. When vector mesons are included, the result is close to the experiments when $Q^2$ is less than 0.4 GeV$^2$ [14]. An alternative regularization method, namely finite-range-regularization (FRR) has been proposed. Inspired by quark models that account for the finite-size of the nucleon as the source of the pion cloud, effective field theory with FRR has been widely applied to extrapolate lattice data of vector meson mass, magnetic moments, magnetic form factors, strange form factors, charge radii, first moments of GPDs, nucleon spin, etc [15–24].

Recently, we proposed a nonlocal chiral effective Lagrangian which makes it possible to study the hadron properties at relatively large $Q^2$ [25, 26]. The nonlocal interaction generates both the regulator which makes the loop integral convergent and the $Q^2$ dependence of form factors at tree level. The obtained electromagnetic form factors and strange form factors of nucleon are very close to the experimental data [25, 26]. This nonlocal chiral effective theory was also applied to study the parton distribution functions and Sivers functions of sea quark in nucleon [27, 28]. In addition, the nonlocal behavior is further assumed to be a general property for all the interactions and an example of this assumption is the application to the lepton anomalous magnetic moments [29].

Since the nonlocal effective theory provides good descriptions to the nucleon form factors up to relatively large momentum transfer. In this paper, we will extend our study from nucleon to all the octet baryons. While the nucleon form factors are precisely determined experimentally, those of the other octet baryons are significantly more challenging to measure and as a result are poorly known from nature. A lot of theoretical and experimental efforts have been made in this field. On the one hand, with the upgrade of the experimental facility, the parton distribution functions (PDFs) from the deep inelastic scattering as well as the form factors at relatively large momentum transfer from the elastic scattering can be extracted [1, 2]. On the other hand, many measurements on form factors have been carried out at very small momentum transfer to get the information of nucleon radii as accurate as possible [3, 4].

In section II, we will introduce the nonlocal chiral Lagrangian. The octet form factors are derived in section III. The paper is organized as follows. In section II, we will introduce the nonlocal chiral Lagrangian. The octet form factors are derived in section III. Numerical results are presented in section IV and finally, section V is a short summary.

II. CHIRAL EFFECTIVE LAGRANGIAN

The lowest order chiral Lagrangian for baryons, pseudo-scalar mesons and their interactions can be written as

\begin{equation}
\mathcal{L} = i Tr \, B_\gamma \, \gamma^\mu B - m_B \, Tr \, BB + \bar{T}^{abc} (i \gamma^{\mu\nu} D_\alpha - m_T \gamma^{\mu\nu}) T_{\nu}^{abc} + \frac{f^2}{4} Tr \, \partial_\mu \Sigma \, \partial^\mu \Sigma^\dagger + D_T \, B_\gamma \, \gamma^5 \{ A_\mu, B \} + F_T \, \bar{B}_\gamma \, \gamma^5 \{ A_\mu, B \} + \frac{C}{f} \epsilon^{abc} \bar{T}_{\mu, a} \, \partial_\mu (g^{\mu\nu} + z_\gamma \gamma_\nu) B_{\nu, a} \phi_{bd} + H.C.,
\end{equation}

where $D$, $F$ and $C$ are the coupling constants. The chiral covariant derivative $D_\mu$ is defined as $D_\mu = \partial_\mu + [V_\mu, B]$. The pseudo-scalar meson octet couples to the baryon field via the vector and axial-vector combinations as

\begin{equation}
V_\mu = \frac{1}{2} (\zeta \partial_\mu \zeta^\dagger + \zeta^\dagger \partial_\mu \zeta) + \frac{1}{2} e \gamma^\mu (\zeta^\dagger Q_c \zeta + \zeta Q_c \zeta^\dagger), \quad A_\mu = \frac{i}{2} (\zeta \partial_\mu \zeta^\dagger - \zeta^\dagger \partial_\mu \zeta) - \frac{1}{2} e \gamma^\mu (\zeta Q_c \zeta^\dagger - \zeta^\dagger Q_c \zeta),
\end{equation}

where $Q_c$ is the charge of the octet baryons.
where
\[ \zeta^2 = \Sigma = e^{i \phi / f}, \quad f = 93 \text{ MeV}. \] (3)

\( Q_c \) is the real charge matrix diag(2/3, -1/3, -1/3). \( \phi \) and \( B \) are the matrices of pseudo-scalar fields and octet baryons. \( \phi^\mu \) is the photon field. The covariant derivative \( D_\mu \) in the decuplet sector is defined as \( D_\mu T_{\mu\nu}^{abc} = \partial_\mu T_{\mu\nu}^{abc} + (\Gamma_\nu, T_{\mu\nu}^{abc}) \), where \( \Gamma_\nu \) is the chiral connection defined as \( (X, T_\mu)^{abc} = (X)_g^T \bar{T}_{\mu}^{abc} + (X)_h^T \bar{T}_{\mu}^{abc} + (X)_i^T \bar{T}_{\mu}^{abc} \). \( \gamma^{\mu\nu}, \gamma^{\mu\nu} \) are the antisymmetric matrices expressed as
\[ \gamma^{\mu\nu} = \frac{1}{2} [\gamma^\mu, \gamma^\nu] \quad \text{and} \quad \gamma^{\mu\nu\rho} = \frac{1}{4} \{[\gamma^\mu, \gamma^\nu], \gamma^\rho\}. \] (4)

The octet, decuplet and octet-decuplet transition operators for magnetic moment are needed in the one loop calculations. The baryon octet anomalous magnetic Lagrangian is written as
\[ \mathcal{L}_{\text{oct}} = \frac{e}{4m_B} \left( c_1 \text{ Tr} [B \sigma^{\mu\nu} \{ F^{+}_{\mu\nu}, B \}] + c_2 \text{ Tr} [B \bar{o}^{\mu\nu} \{ F^{+}_{\mu\nu}, B \}] + c_3 \text{ Tr} [B \bar{o}^{\mu\nu} B] \text{ Tr} [F^{+}_{\mu\nu}, B] \right), \] (5)
where,
\[ F^{+}_{\mu\nu} = -\frac{1}{2} \left( \zeta^\dagger F_{\mu\nu} Q_c \zeta + \zeta F_{\mu\nu} Q_c \zeta^\dagger \right). \] (6)

At the lowest order, the contribution of quark \( q \) with unit charge to the octet magnetic moments can be obtained by replacing the charge matrix \( Q_c \) with the corresponding diagonal quark matrices \( \lambda_q = \text{diag}(\delta_{qu}, \delta_{qd}, \delta_{qs}) \). Let’s take the nucleon as an example. After expansion of the above equation, it is found that
\[ F^{p,u}_2 = c_1 + c_2 + c_3, \quad F^{p,d}_2 = c_4, \quad F^{p,s}_2 = c_1 - c_2 + c_3, \quad F^{n,u}_2 = c_3, \quad F^{n,d}_2 = c_1 + c_2 + c_3, \quad F^{n,s}_2 = c_1 - c_2 + c_3. \] (7)

Comparing with the results of the constituent quark model where
\[ F^{p,s}_2 = 0 \quad \text{and} \quad F^{n,s}_2 = 0, \] (8)
we get
\[ c_3 = c_2 - c_1. \] (9)

The decuplet anomalous magnetic moment operator is expressed as
\[ \mathcal{L}_{\text{dec}} = -\frac{ieF^T_2}{4M_T} \bar{T}_{\mu,abc} \sigma_\rho \lambda T^{\mu,abc}. \] (10)

The transition magnetic operator is written as
\[ \mathcal{L}_{\text{trans}} = \frac{e}{4m_B} \mu_T F_{\mu\nu} \left( \epsilon^{ijk} Q_{c,il} \bar{B}_{jm} \gamma^\mu \gamma_5 T^{\nu,klm} + \epsilon^{ijk} Q_{c,il} \bar{T}_{\mu,klm} \gamma^\nu \gamma_5 B_{mj} \right). \] (11)

The anomalous magnetic moments of baryons can also be expressed in terms of quark magnetic moments \( \mu_q \). For example, \( \mu_p = \frac{2}{5} \mu_u - \frac{3}{5} \mu_d, \mu_n = \frac{4}{5} \mu_d - \frac{3}{5} \mu_u, \mu_{\Delta^{++}} = 3 \mu_u \). Using the SU(3) symmetry, \( \mu_u = -2 \mu_d = -2 \mu_s, \mu_T \) and \( F^T_2 \) as well as \( \mu_q \) can be written in terms of \( c_1 \) or \( c_2 \). For example, \( \mu_u = \frac{2}{3} c_1, \mu_T = 4 c_1, F^{\Delta^{++}}_2 = \mu_{\Delta^{++}} - 2 = 2c_1 - 2. \)

The gauge invariant nonlocal Lagrangian can be obtained using the method in [25, 26, 35]. For instance, the local interaction between hyperons and \( K^- \) meson can be written as
\[ \mathcal{L}^{\text{local}}_K = \int dx \frac{D + F}{\sqrt{2f}} \Sigma^0(x) \gamma^\mu \gamma_5 (x)(\partial_\mu + ie \mathcal{A}_\mu(x)) K^-(x). \] (12)
The corresponding nonlocal Lagrangian is expressed as

\[
\mathcal{L}_{\text{nl}}^\text{local} = -e \int da \bar{\Sigma}^+(x) \gamma^\mu \Sigma^+(x) \sigma_{\mu
u} \Sigma(x)^+ F_{\mu\nu}(x).
\] (14)

The corresponding nonlocal Lagrangian is expressed as

\[
\mathcal{L}_{\text{EM}}^\text{nl} = -e \int da \bar{\Sigma}^+(x) \gamma^\mu \Sigma(x)^+ \sigma_{\mu
u} \Sigma(x)^+ F_{\mu\nu}(x-a) F_1(a) + \frac{(c_1 + 3c_2)e}{12m_{\Sigma}} \int da \bar{\Sigma}^+(x) \sigma_{\mu
u} \Sigma(x)^+ F_{\mu\nu}(x-a) F_2(a),
\] (15)

where \(F_1(a)\) and \(F_2(a)\) are the correlation functions for the nonlocal electric and magnetic interactions.

The momentum dependence of the form factors at tree level can be easily obtained with the Fourier transformation of the correlation function. As in our previous work \([25, 26]\), the correlation function is chosen such that the charge conservation is guaranteed.

The nonlocal electromagnetic interaction can also be obtained in the same procedure. For example, the local interaction between \(\Sigma^+\) and photon is written as

\[
\int dy \frac{D + F}{\sqrt{2f}} \Xi^0(x) \gamma^\mu \Sigma^+(x) \sigma_{\mu
u} \Sigma(x)^+ F_{\mu\nu}(x),
\] (13)

where \(F(x)\) is the correlation function. To guarantee gauge invariance, the gauge link is introduced in the above Lagrangian. With the correlation function, the regulator can be generated automatically.

The electromagnetic form-factors are defined as the combinations of the above form factors as

\[
G_E^B(Q^2) = F_1^B(Q^2) - \frac{Q^2}{4m_B^2} F_2^B(Q^2), \quad G_M^B(Q^2) = F_1^B(Q^2) + F_2^B(Q^2).
\] (20)
With the electromagnetic form factors, the magnetic and electric (charge) radii can be obtained. The magnetic radii of octet baryons are defined as

$$\langle r_M^2 \rangle_B = -\frac{6}{G_M^B(0)} \frac{dG_M^B(Q^2)}{dQ^2} \mid_{Q^2=0}. \quad (21)$$

The electric radii of charged and neutral baryons are defined as

$$\langle r_E^2 \rangle_B = -\frac{6}{G_E^B(0)} \frac{dG_E^B(Q^2)}{dQ^2} \mid_{Q^2=0} \text{ and } \langle r_E^2 \rangle_B = -\frac{6}{dG_E^B(Q^2)} \mid_{Q^2=0}, \quad (22)$$

respectively.

According to the Lagrangian, the one loop Feynman diagrams which contribute to the octet electromagnetic form factors are shown in Fig. 1. From the Lagrangian, we can get the matrix element of Eq. (19). Let’s take the Σ hyperons as an example. In this section, we only show the expressions for the intermediate octet baryons. For the intermediate decuplet baryons, the expressions are similar but more complicated which are expressed in Appendix A.

The contributions of Fig. 1a are written as

$$\Gamma_{\mu}^\phi(\Sigma^-) = \frac{F_\phi^2}{f_\phi^2} I_a^{\Sigma\pi} + \frac{D_\phi^2}{3f_\phi^2} I_a^{\Lambda\pi} + \frac{(D - F_\phi)^2}{2f_\phi^2} I_a^{N,K}, \quad (23)$$

$$\Gamma_{\mu}^\phi(\Sigma^0) = \frac{(D - F_\phi)^2}{4f_\phi^2} I_a^{N,K} - \frac{(D + F_\phi)^2}{4f_\phi^2} I_a^{\Xi,K}, \quad (24)$$

$$\Gamma_{\mu}^\phi(\Sigma^+) = -\frac{F_\phi^2}{f_\phi^2} I_a^{\Sigma\pi} - \frac{D_\phi^2}{3f_\phi^2} I_a^{\Lambda\pi} - \frac{(D + F_\phi)^2}{2f_\phi^2} I_a^{\Xi,K}, \quad (25)$$

where the integral $I_a^{BM}$ is expressed as

$$I_a^{BM} = \bar{u}(p') \hat{F}(q) \int \frac{d^4k}{(2\pi)^4} \frac{\hat{F}(q + k)\hat{F}(k)}{D_M(k + q)} - (2k + q)^{\mu} \frac{1}{\not{k} - \not{p} - m_B} \not{k} \gamma_5 u(p). \quad (26)$$
where \(m\) is defined as

\[
D_M(k) = k^2 - m_M^2 + i\epsilon. \tag{27}
\]

\(m_B\) and \(m_M\) are the masses of the intermediate B baryon and \(M\) meson.

The contributions of Fig. 1b are expressed as

\[
\Gamma_\mu^\nu(\Sigma^-) = \Gamma_\mu^\nu(\Sigma^0) = \frac{2c_1 F^2 Q_2}{3f^2 (4m_N^2 + Q^2)} I_\nu^K - \frac{c_1 D Q_2 (D - 3F)}{9f^2 (4m_A^2 + Q^2)} I_\nu^K,
\]

\[
\Gamma_\mu^\nu(\Sigma^+) = \frac{2c_1 F^2 Q_2}{3f^2 (4m_N^2 + Q^2)} I_\nu^K - \frac{c_1 D Q_2 (D + 3F)}{9f^2 (4m_A^2 + Q^2)} I_\nu^K,
\]

\[
\Gamma_\mu^\nu(\Sigma^+) = \frac{2c_1 F^2 Q_2}{3f^2 (4m_N^2 + Q^2)} I_\nu^K - \frac{c_1 D Q_2 (D + 3F)}{9f^2 (4m_A^2 + Q^2)} I_\nu^K,
\]

where the integral \(I_{BM}^B\) is written as

\[
I_{BM}^B = \bar{u}(p') \tilde{F}(q) \int \frac{d^4 k}{(2\pi)^4} \frac{\tilde{F}(k)^2}{D_M(k)} \frac{1}{\not{k} - \not{p} - m_B} \gamma^\mu \frac{1}{\not{k} - \not{p} - m_B} \tilde{\gamma}_5 u(p). \tag{31}
\]

Fig. 1c is similar to Fig. 1b except for the magnetic interaction. The contributions of this diagram are written as

\[
\Gamma_\mu^\nu(\Sigma^-) = \frac{2F m_\Sigma (c_1 (D + 2F) - 3c_2 F)}{3f^2 (4m_N^2 + Q^2)} I_{e^K} - \frac{2c_1 D (D - 3F) m_\Lambda}{9f^2 (4m_A^2 + Q^2)} I_{e^K},
\]

\[
\Gamma_\mu^\nu(\Sigma^0) = \frac{2c_1 F^2 m_\Sigma}{3f^2 (4m_B^2 + Q^2)} I_{e^K} - \frac{2c_1 D^2 m_\Lambda}{9f^2 (4m_A^2 + Q^2)} I_{e^K},
\]

\[
\Gamma_\mu^\nu(\Sigma^+) = \frac{2F m_\Sigma (c_1 (D) + 2c_1 F + 3c_2 F)}{3f^2 (4m_B^2 + Q^2)} I_{e^K} - \frac{2c_1 D (D + 3F) m_\Lambda}{9f^2 (4m_A^2 + Q^2)} I_{e^K},
\]

where \(I_{eBM}^B\) is expressed as

\[
I_{eBM}^B = \bar{u}(p') \tilde{F}(q) \int \frac{d^4 k}{(2\pi)^4} \frac{\tilde{F}(k)^2}{D_M(k)} \frac{1}{\not{k} - \not{p} - m_B} i\sigma^{\mu\nu} q_{\nu} \frac{1}{\not{k} - \not{p} - m_B} \tilde{\gamma}_5 u(p). \tag{35}
\]
The coupling constants between octet baryons and mesons are determined by the two parameters \( \mu \) chosen as 2. The off-shell parameter \( z \) is chosen to be -1. The covariant derivative in Package-X [36] is used to simplify the loop integral. The expressions for the Dirac and Pauli form factors are derived. In the next section, the numerical results are discussed.

### IV. Numerical Results

The coupling constants between octet baryons and mesons are determined by the two parameters \( D \) and \( F \). The numerical calculations, the parameters are chosen as \( D = 0.76 \) and \( F = 0.50 \). The coupling constant \( C \) is chosen to be 2 which is the same as in Ref. [25, 41]. The numerical results are presented in Table I.

| \( \mu \)  | Tree | Loop | Total | Lattice [37] | Lattice [31] | NJL [38] | PCQM [39] | Exp. [40] |
|----------|------|------|-------|------------|------------|-------|---------|---------|
| \( \mu_\pi \) | 1.850 | 0.795 | 2.644 ± 0.159 | 2.4(2) | 2.3(3) | 2.78 | 2.735 ± 0.121 | 2.793 |
| \( \mu_n \) | -0.859 | -1.126 | -1.984 ± 0.216 | -1.59(17) | -1.45(17) | -1.81 | -1.956 ± 0.103 | -1.913 |
| \( \mu_{\Sigma^+} \) | 1.850 | 0.572 | 2.421 ± 0.147 | 2.27(16) | 2.12(18) | 2.62 | 2.537 ± 0.201 | 2.458 ± 0.010 |
| \( \mu_{\Sigma^0} \) | 0.429 | 0.155 | 0.584 ± 0.057 | - | - | - | 0.838 ± 0.091 | - |
| \( \mu_{\Sigma^-} \) | -0.991 | -0.262 | -0.594 ± 0.057 | - | - | - | 0.861 ± 0.040 | 0.610 ± 0.025 |
| \( \mu_{\Xi^0} \) | -0.859 | -0.521 | -1.380 ± 0.169 | -1.32(4) | -1.07(7) | -1.14 | -1.690 ± 0.142 | -1.250 ± 0.014 |
| \( \mu_{\Xi^-} \) | -0.991 | 0.266 | -0.725 ± 0.077 | -0.71(3) | -0.57(5) | -0.67 | -0.840 ± 0.087 | -0.651 ± 0.080 |

Table I. The tree, loop and total contributions to the octet magnetic moments \( \mu_B \) (in units of the nucleon magneton \( \mu_N \)). The results from two lattice simulations, NJL and PCQM models as well as the experimental data are also listed.
The regulator is chosen to be dipole form [25–27]

$$\tilde{F}[k] = \frac{\Lambda^4}{(k^2 - m_j^2 - \Lambda^2)^2},$$

(44)

where $m_j$ is the meson mass for the baryon-meson interaction and it is zero for the hadron-photon interaction. It was found that when $\Lambda$ was around 0.90 GeV, the obtained nucleon form factors are very close to the experimental data. Therefore, all the above parameters are predetermined. There are only two free parameters which are the low energy constants $c_1$ and $c_2$. In our previous calculation for the nucleon form factors, they are fitted by the experimental nucleon magnetic moments. Here, $c_1$ and $c_2$ are determined to be 1.288 and 0.420, which give the minimal standard variation $\chi^2$ of the octet magnetic moments.

In Table I, the tree, loop and total contributions to the baryon magnetic moments obtained from the nonlocal chiral effective theory are listed. The error bar in our calculation is determined by varying $\Lambda$ from 0.8 to 1.0 GeV. The results from two lattice simulations, NJL and PCQM models as well as the experimental data are also listed for comparison. From the table, one can see that all the magnetic moments of baryon octets are reasonably reproduced. The largest deviation from the experiments is for $\Xi_s$, where the calculated absolute values of magnetic moments of $\mu_{\Xi^0}$ and $\mu_{\Xi^-}$ are about 10% larger than experimental ones. For the other baryons, the deviation from the experiments is less than 5%. It is interesting that all the tree and loop contributions have the same signs except for $\Xi^-$, where the loop diagrams give the opposite contribution to the tree diagram. The data from lattice simulations are somewhat smaller which is partially due to the large pion mass and/or the neglecting of the disconnected contribution.

The magnetic form factors of charged octet baryons versus momentum transfer $Q^2$ are plotted in Fig. 2. The solid, dashed, dotted and dash-dotted lines are for proton, $\Sigma^+$, $\Sigma^-$ and $\Xi^-$, respectively. It is clear the proton magnetic form factor is comparable with the experimental data up to 1 GeV$^2$. This is the advantage of the nonlocal chiral effective theory. The correlation function in the nonlocal Lagrangian makes the loop integral convergent. In the mean time, it provides the momentum dependence of the form factors at tree level, and as a result, the total form factors can be close to the experimental data up to relatively large $Q^2$. The other charged form factors have the similar momentum dependence as proton. Among them, $\Xi^-$ decreases a little slower with the increasing $Q^2$. The magnetic radii are determined by the slopes of the form factors at zero momentum transfer which will be discussed later.

The normalized magnetic form factors for the charge neutral baryons are plotted in Fig. 3. The solid, dashed, dotted and dash-dotted lines are for neutron, $\Sigma^0$, $\Lambda$ and $\Xi^0$, respectively. The magnetic from factors of $\Sigma^0$, $\Lambda$ and $\Xi^0$ are close to each other. The neutron magnetic form factor is a little smaller than the experimental data and it drops faster than the other three neutral baryons. All of the form factors have a dipole-like momentum dependence.

The tree, loop and total contributions to the magnetic radii of octet baryons are listed in Table II. The data from lattice simulation and phenomenological quark models are also listed for comparison. Though our central values for
proton and neutron are a little larger than experiments, the results are still reasonable. The magnetic radii of octet baryons vary from 0.5 fm$^2$ to 0.9 fm$^2$, but show no simple dependence on baryon/quark mass. $\Sigma^-$ has the largest contribution at tree level. Because of the opposite contribution from the loop diagrams, its total radius is the smallest among all the neutral baryons. The experimental form factor of neutron is from Refs. [46–57].

We now discuss the electric form factors. In Fig. 4, we plot the electric form factors of the charged baryons. Because of the additional interaction which makes the nonlocal Lagrangian locally gauge invariant, the electric form factors start from their charge at $Q^2 = 0$. The proton charge form factors is close to the experimental data. The absolute values of the electric form factors have the similar momentum dependence. This could be examined by the further experiments and/or accurate lattice simulation.

The electric form factors for the neutral baryons are plotted in Fig. 5. Again due to the charge conservation, the form factors start from 0 at zero momentum transfer. The calculated electric form factor of neutron is consistent with the experimental data. The form factors of the other neutral baryons are very small. There is no tree level contribution to the electric form factors of neutral baryons and all the contributions are from the loop diagrams. Among them, neutron has the largest contribution from $\pi$-loop diagrams. The corresponding $\pi$-loop diagrams for the other neutral baryons are fairly small due to the small coupling constants.

The charge radii of octet baryons are listed in Table III. Our results are comparable with the experimental data in PDG for nucleon and $\Sigma^-$. A small proton charge radius $\langle r_E^2 \rangle_p = 0.831 \pm 0.067 \pm 0.012$, i.e. $\langle r_{E,\Sigma}^2 \rangle_p = 0.691 \pm 0.032$.
was reported recently [3] which is also close to our value $\langle r_E^2 \rangle_p = 0.729 \pm 0.112$. For the neutral baryons, the loop contribution is very small except neutron. For the charged baryons, the tree level contributions are the same which are also dominant for all of them. The loop contribution has the same order of magnitude except for $\Xi^-$, where the loop contribution is small. Different from the magnetic radii, the total charge radii vary around 0.6 and 0.7 fm$^2$ for the charged baryons.

FIG. 4. Same as Fig. 2 but for electric form factors, The experimental form factor of proton is from Refs. [43–45, 47, 48, 58–61].

FIG. 5. Same as Fig. 3 but for electric form factors, The experimental form factor of neutron is from Refs. [50, 52, 56, 62–68].

V. SUMMARY

We applied the nonlocal chiral effective theory to study the electromagnetic form factors of octet baryons. The correlation function in the Lagrangian makes the loop integral convergent. It also provides the momentum dependence of the form factors at tree level. The additional interaction generated from the expansion of the gauge link guarantees the Lagrangian is locally gauge invariant. This nonlocal Lagrangian makes it possible to study the physical quantities at relatively large momentum transfer in the framework of chiral effective theory. In the numerical calculation, all the parameters are predetermined except the two low energy constants $c_1$ an $c_2$. They are fitted to give the minimum
of the standard deviation. When extending the previous study of form factors of nucleon to all the baryon octets, we do not add any new parameter. The magnetic moments are well reproduced. The deviation from the experiments is less than 5% except \(\Xi^0\) and \(\Xi^−\), where the deviation is about 10%. For the radii, most experiments focus on the nucleon and there is few data for the other baryons. Our results on magnetic and charge radii are comparable with those of other theoretical methods.

For the other octet baryons, since the method is the same, we expect this nonlocal Lagrangian can also give good descriptions. The difference between our results and those of other theoretical methods could be examined by future experiments and more accurate lattice simulations.

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Appendix A: Decuplet Expressions

In this section, we show the expressions of one loop integrals for decuplet intermediate states. The contribution for Fig. 1h can be written as

\[
\Gamma_h^\mu(\Sigma^-) = \frac{C^2}{12f^2} I_h^{\Sigma^\gamma\pi} + \frac{C^2}{6f^2} I_h^{\Delta K},
\]

(A1)

\[
\Gamma_h^\mu(\Sigma^0) = \frac{C^2}{12f^2} I_h^{\Sigma^\gamma\pi} - \frac{C^2}{3f^2} I_h^{\Delta K},
\]

(A2)

\[
\Gamma_h^\mu(\Sigma^+) = \frac{C^2}{12f^2} I_h^{\Sigma^\gamma\pi} - \frac{C^2}{6f^2} I_h^{\Delta K} + \frac{C^2}{2f^2} I_h^{\Delta K},
\]

(A3)

where the integral \(I_h^{TM}\) is expressed as

\[
I_h^{TM} = \bar{u}(p') \hat{F}(q) \int \frac{d^4k}{(2\pi)^4} \frac{\hat{F}(q+k)\hat{F}(k)}{D_M(k)} \frac{2(k+q)\mu}{\not{p} - \not{k} - m_T} \cdot S_{\sigma\rho}(p-k)(-k^\rho - z\gamma^\rho\gamma^5)u(p).
\]

(A4)

\[
\times((k+q)\sigma + z(\not{k} + \not{q})\gamma^\sigma) \frac{1}{\not{p} - \not{k} - m_T} \cdot S_{\sigma\rho}(p-k)(-k^\rho - z\gamma^\rho\gamma^5)u(p).
\]

(A5)

\(m_T\) is the mass of the decuplet intermediate state and \(S_{\sigma\rho}(p)\) is expressed as

\[
S_{\sigma\rho}(p) = -g_{\sigma\rho} + \frac{\gamma_\sigma \gamma_\rho}{3} + \frac{p_\sigma p_\rho}{3m_T^2} + \frac{\gamma_\sigma p_\rho - \gamma_\rho p_\sigma}{3m_T}.
\]

(A6)
The contribution for Fig. 1i is written as

\[ \Gamma_i^\mu (\Sigma^-) = - \frac{C^2 ((c_1 + 3c_2 + 3) Q^2 + 12m_{\Sigma^*}^2)(I_i^{\Sigma^-})}{36f^2 (4m_{\Sigma^*}^2 + Q^2)} - \frac{C^2 ((c_1 + 3c_2 + 3) Q^2 + 12m_{\Sigma^*}^2)(I_i^{\Sigma^+})}{18f^2 (4m_{\Sigma^*}^2 + Q^2)} \]

\[ + \frac{C^2 ((c_1 + 3c_2 + 3) Q^2 + 12m_{\Delta}^2)(I_i^{\Delta K})}{6f^2 (4m_{\Delta}^2 + Q^2)} \]

\[ - \frac{C^2 ((c_1 + 3c_2 + 3) Q^2 + 12m_{\Delta}^2)(I_i^{\pi K})}{36f^2 (4m_{\Delta}^2 + Q^2)} \]

\[ - \frac{C^2 ((c_1 + 3c_2 + 3) Q^2 + 12m_{\Delta}^2)(I_i^{\pi K})}{18f^2 (4m_{\Delta}^2 + Q^2)} \]

\[ + \frac{7C^2 ((c_1 + 3c_2 + 3) Q^2 + 12m_{\Delta}^2)(I_i^{\Delta K})}{18f^2 (4m_{\Delta}^2 + Q^2)} \]

where the integral \( I_i^{TM} \) is written as

\[ I_i^{TM} = \bar{u}(p') \hat{F}(q) \int \frac{d^4k}{(2\pi)^4} \frac{\hat{F}(k)^2}{D_M(k)} (k^\sigma + zk^\sigma) \]

\[ \times \frac{1}{p' - k - m_T} S_{\sigma\nu}(p' - k) \gamma^\alpha \gamma^\lambda q_{\lambda\nu} \frac{1}{p' - k - m_T} S_{\nu\rho}(p - k)(k^\rho + z\gamma^\rho \bar{k} u(p). \]

The contribution for Fig. 1j is written as

\[ \Gamma_j^\mu (\Sigma^-) = - \frac{(c_1 + 3c_2) C^2 m_{\Sigma^*}}{18f^2 (4m_{\Sigma^*}^2 + Q^2)} \frac{I_j^{\Sigma^-}}{I_j^{\Sigma^+}} + \frac{(c_1 + 3c_2) C^2 m_{\Delta}}{3f^2 (4m_{\Delta}^2 + Q^2)} \frac{I_j^{\Delta K}}{I_j^{\Delta K}} + \frac{(c_1 + 3c_2) C^2 m_{\Delta}}{9f^2 (4m_{\Delta}^2 + Q^2)} \frac{I_j^{\pi K}}{I_j^{\pi K}} \]

\[ - \frac{(c_1 + 3c_2) C^2 m_{\Delta}}{18f^2 (4m_{\Delta}^2 + Q^2)} \frac{I_j^{\Delta K}}{I_j^{\Delta K}} + \frac{(c_1 + 3c_2) C^2 m_{\Delta}}{9f^2 (4m_{\Delta}^2 + Q^2)} \frac{I_j^{\pi K}}{I_j^{\pi K}} \]

where the integral \( I_j^{TM} \) is expressed as

\[ I_j^{TM} = \bar{u}(p') \hat{F}(q) \int \frac{d^4k}{(2\pi)^4} \frac{\hat{F}(k)^2}{D_M(k)} (k^\sigma + zk^\sigma) \]

\[ \times \frac{1}{p' - k - m_T} S_{\sigma\nu}(p' - k) i\sigma^\mu \lambda q_{\lambda\nu} \frac{1}{p' - k - m_T} S_{\nu\rho}(p - k)(k^\rho + z\gamma^\rho \bar{k} u(p). \]

The contribution for the intermediate octet-decuplet transition diagrams Fig. 1k & l is expressed as

\[ \Gamma_{k+1}^\mu (\Sigma^-) = - \frac{c_1 CF}{12f^2 m_{\Sigma}} I_{k+1}^{\Sigma^-} + \frac{c_1 CD}{12f^2 m_{\Delta}} I_{k+1}^{\Sigma^-} \]

\[ + \frac{c_1 C(D - F)}{6f^2 m_{N}} \frac{I_{k+1}^{\Delta K}}{I_{k+1}^{\Delta K}} \]

\[ - \frac{c_1 C(D - F)}{12f^2 m_{\Delta}} \frac{I_{k+1}^{\pi K}}{I_{k+1}^{\pi K}} \]

where the integral \( I_{k+1}^{TM} \) is written as,

\[ I_{k+1}^{TM} = \bar{u}(p') \hat{F}(q) \int \frac{d^4k}{(2\pi)^4} \frac{\hat{F}(k)^2}{D_M(k)} \frac{1}{p' - k - m_B} \gamma^\mu \gamma^\nu q_{\nu\mu} \frac{1}{p' - k - m_T} S_{\nu\rho}(p' - k)(k^\rho + z\gamma^\rho \bar{k} u(p) \]

\[ + (k^\sigma + z\gamma^\sigma) \frac{1}{p' - k - m_T} S_{\nu\rho}(p' - k) q_{\nu\mu} \frac{1}{p' - k - m_B} \frac{\hat{F}(k)^2}{D_M(k)} \]

\[ + (k^\sigma + z\gamma^\sigma) \frac{1}{p' - k - m_T} S_{\nu\rho}(p' - k) q_{\nu\mu} \frac{1}{p' - k - m_B} \frac{\hat{F}(k)^2}{D_M(k)} \]

\[ + (k^\sigma + z\gamma^\sigma) \frac{1}{p' - k - m_T} S_{\nu\rho}(p' - k) q_{\nu\mu} \frac{1}{p' - k - m_B} \frac{\hat{F}(k)^2}{D_M(k)} \]

\[ + (k^\sigma + z\gamma^\sigma) \frac{1}{p' - k - m_T} S_{\nu\rho}(p' - k) q_{\nu\mu} \frac{1}{p' - k - m_B} \frac{\hat{F}(k)^2}{D_M(k)} \]

\[ \bar{u}(p') \hat{F}(q) \int \frac{d^4k}{(2\pi)^4} \frac{\hat{F}(k)^2}{D_M(k)} \frac{1}{p' - k - m_B} \gamma^\mu \gamma^\nu q_{\nu\mu} \frac{1}{p' - k - m_T} S_{\nu\rho}(p' - k)(k^\rho + z\gamma^\rho \bar{k} u(p). \]
The contribution for the Kroll-Ruderman diagrams Fig. 1 m & n is written as
\[
\Gamma_{m+n}^{\mu}(\Sigma^-) = \frac{C^2}{12f^2} I_{m+n}^{\Sigma^-} + \frac{C^2}{6f^2} I_{m+n}^{\Delta K},
\]
(A21)
\[
\Gamma_{m+n}^{\mu}(\Sigma^0) = \frac{C^2}{3f^2} I_{m+n}^{\Delta K} - \frac{C^2}{12f^2} I_{m+n}^{\Sigma^+ K},
\]
(A22)
\[
\Gamma_{m+n}^{\mu}(\Sigma^+) = -\frac{C^2}{12f^2} I_{m+n}^{\Sigma^+} + \frac{C^2}{2f^2} I_{m+n}^{\Delta K} - \frac{C^2}{6f^2} I_{m+n}^{\Sigma^+ K},
\]
(A23)
where the integral \( I_{m+n}^{TM} \) is written as,
\[
I_{m+n}^{TM} = \bar{u}(p') \hat{F}(q) \int \frac{d^4k}{(2\pi)^4} \frac{\hat{F}(k)^2}{M_D(k)} \left\{ (k + z\Sigma^-) \frac{1}{p' - \gamma^2 - m_T} S^{\rho}(p' - k) (g_{\rho}\mu + z\gamma^\rho \gamma^\mu) + (g_{\rho}\sigma + z\gamma^\rho \gamma^\sigma) \frac{1}{p' - \gamma^2 - m_T} S^{\rho}(p - k) (k_{\rho} + z\gamma_{\rho} \gamma_{\rho}) \right\} u(p).
\]
(A24)

The contribution for the additional diagrams with intermediate decuplet states Fig. 1 o & p is expressed as
\[
\Gamma_{o+p}^{\mu}(\Sigma^-) = \frac{C^2}{12f^2} I_{o+p}^{\Sigma^-} + \frac{C^2}{6f^2} I_{o+p}^{\Delta K},
\]
(A25)
\[
\Gamma_{o+p}^{\mu}(\Sigma^0) = \frac{C^2}{3f^2} I_{o+p}^{\Delta K} - \frac{C^2}{12f^2} I_{o+p}^{\Sigma^+ K},
\]
(A26)
\[
\Gamma_{o+p}^{\mu}(\Sigma^+) = -\frac{C^2}{12f^2} I_{o+p}^{\Sigma^+} + \frac{C^2}{2f^2} I_{o+p}^{\Delta K} - \frac{C^2}{6f^2} I_{o+p}^{\Sigma^+ K},
\]
(A27)
where the integral \( I_{o+p}^{TM} \) is written as
\[
I_{o+p}^{TM} = \bar{u}(p') \hat{F}(q) \int \frac{d^4k}{(2\pi)^4} \frac{\hat{F}(k)}{M_D(k)} \left\{ \frac{(2k+q)^{\mu}}{2k + q^2} \left( \hat{F}(k - q) - \hat{F}(k) \right) \frac{1}{p' - \gamma^2 - m_T} S^{\rho}(p' - k) (k_{\rho} + z\gamma_{\rho} \gamma_{\rho}) \right\} u(p).
\]
(A28)

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