Electric Polarizabilities of Neutral Baryons in the QCD Sum Rule

Tetsuo NISHIKAWA and Sakae SAI TO

Department of Physics, Nagoya University, Nagoya 464-01, Japan

Yoshihiko KONDO

Kokugakuin University, Tokyo 150, Japan

Abstract

We investigate the electric polarizabilities of neutral baryons using the method of QCD sum rules. The diagrams in the operator product expansion are taken into account up to dimension 6. We obtain different values of the polarizabilities among hyperons, the ordering of which is given $\bar{\alpha}_n > \bar{\alpha}_\Xi^0 > \bar{\alpha}_\Lambda$. The polarizability of $\Sigma^0$, $\bar{\alpha}_{\Sigma^0}$, has possibly a very small value.

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1E-mail address nishi@nuc-th.phys.nagoya-u.ac.jp
2E-mail address saito@nuc-th.phys.nagoya-u.ac.jp
3E-mail address kondo@kokugakuin.ac.jp
Electric and magnetic polarizabilities, labeled $\bar{\alpha}$ and $\bar{\beta}$, are fundamental constants characterizing the hadron structure. These constants measure the ease with which a hadron acquires electric and magnetic dipole moments in external electromagnetic fields. While for the proton and neutron they have long been studied through theoretically and experimentally and recent measurements have established certain values, very little is known about hyperon polarizabilities. However with the hyperon beams at CERN and Fermilab, the situation is expected to change. In particular the $\Sigma$ polarizabilities will be soon measured.

This has induced a number of theoretical investigations. In fact predictions have been made in the non-relativistic quark model, the heavy baryon chiral perturbation theory and the bound state soliton model. In the non-relativistic quark model the results are $\bar{\alpha}_{\Sigma^+} = 20.8$, $\bar{\alpha}_{\Sigma^-} = 5.7$, all in units of $10^{-4}\text{fm}^3$. The large difference in them can be traced back to the different quark structure of the $\Sigma^+$ and $\Sigma^-$. While in the $\Sigma^+$ the two quark flavors have opposite electric charges, all valence quarks in the $\Sigma^-$ have the same charges, which inhibit internal excitation. In agreement with the non-relativistic quark model, the heavy baryon chiral perturbation theory (HBChPT) at leading order finds $\bar{\alpha}_{\Sigma^+} > \bar{\alpha}_{\Sigma^-}$; namely, $\bar{\alpha}_{\Sigma^+} = 9.4(8.8)$ and $\bar{\alpha}_{\Sigma^-} = 6.3(5.9)$ in the same units for the axial-vector coupling constants $D = 0.8(0.75)$ and $F = 0.5(0.5)$. Here the reason for this result is the kaon cloud contribution, which is strongly suppressed for the $\Sigma^-$. In addition hyperon polarizabilities are smaller than the nucleon polarizabilities ($\bar{\alpha}_p = 13.0(12.1)$, $\bar{\alpha}_n = 11.4(10.5)$) in the HBChPT. The large value of $\bar{\alpha}_N$ in the HBChPT is due to the large contribution of pion cloud. On the other hand the hyperon polarizabilities mainly come from the contribution of kaon cloud, whose mass is larger than that of pion and inhibit internal excitation under the external field.

We use a completely alternative approach, the method of QCD sum rules. In a previous work we constructed the QCD sum rule for the electromagnetic polarizabilities of the neutron and found that they are related with the vacuum condensates of quark-gluon fields in a weak external constant electromagnetic field. We have seen that the neutron electric polarizability is well described with the method. In this paper we extend the method to hyperons and calculate the electric polarizabilities of neutral hyperons.

We consider a two-point correlation function in a weak external constant electromagnetic field $F_{\mu\nu}$ as follows:

$$\Pi_F(p) = i \int d^4x \, e^{ipx} \langle 0 | T[\eta_B(x)\bar{\eta}_B(0)] | 0 \rangle_F,$$

where $\eta_B(x)$ is an interpolating field constructed out of quark fields such that it has the quantum numbers of a baryon under consideration. Following Ioffe and Smilga, we
expand the correlation function in powers of $F_{\mu \nu}$ up to order $F^2$ terms:

$$\Pi_F(p) = \Pi^{(0)}(p) + \Pi^{(1)}_{\mu \nu}(p)F^\mu \nu + \Pi^{(2)}_{\mu \nu \rho \sigma}(p)F^\mu \nu F^\rho \sigma + \mathcal{O}(F^3). \tag{2}$$

The leading term $\Pi^{(0)}(p)$ is the correlation function when the external field is absent. $\Pi^{(1)}_{\mu \nu}(p)$ is known to be related with the magnetic moments of baryons. It is $\Pi^{(2)}_{\mu \nu \rho \sigma}(p)$ that is related with the electromagnetic polarizabilities of baryons. Requiring parity invariance, we can decompose $\Pi^{(2)}_{\mu \nu \rho \sigma}(p)$ into the following Lorentz structure:

$$\Pi^{(2)}_{\mu \nu \rho \sigma}(p) = g_{\mu \rho}g_{\nu \sigma}\Pi_{S_1}(p^2) + p_\mu p_\rho g_{\nu \sigma}\Pi_{S_2}(p^2) + \sigma_{\mu \lambda}p_\lambda p_{\rho \sigma}\Pi_{T_1}(p^2) + i\gamma_5 \epsilon_{\mu \rho \lambda \sigma}[g_{\lambda \sigma}\Pi_{P_1}(p^2) + p_\lambda p_\sigma\Pi_{P_2}(p^2)]$$

$$+ i\gamma_5 \epsilon_{\mu \rho \lambda \sigma}[g_{\lambda \sigma}p_\rho p_\sigma\Pi_{A_1}(p^2) + \gamma^\rho p_\sigma\Pi_{A_2}(p^2) + \gamma^\rho p_\sigma p_\rho\Pi_{A_3}(p^2) + i\gamma_5 \epsilon_{\mu \rho \lambda \sigma}[g_{\lambda \sigma}p_\rho p_\sigma\Pi_{A_4}(p^2)]. \tag{3}$$

The invariant functions in the right hand side satisfy dispersion relations:

$$\Pi(p^2) = \frac{1}{\pi} \int_0^\infty dt \frac{\text{Im}\Pi(t)}{t - p^2} + \text{subtraction terms}, \tag{4}$$

where the subscripts $S_1$ etc. are suppressed for brevity. One can apply the Borel transform $\hat{B}$ defined by

$$\hat{B}[\Pi(p^2)] \equiv \lim_{s \rightarrow \infty} \lim_{n \rightarrow \infty} \frac{(-p^2)^{n+1}}{n!} \left[ -\frac{d}{d(-p^2)} \right]^n \Pi(p^2). \tag{5}$$

Here $s$ denotes the squared Borel mass. We make a Borel transform on both sides of the dispersion relations:

$$\hat{B}[\Pi(p^2)] = \frac{1}{\pi} \int_0^\infty dt e^{-t/s} \text{Im}\Pi(t). \tag{6}$$

Evaluating the left hand side by an operator product expansion (OPE) in the deep Euclidean region we obtain the Borel sum rules.

Let us next consider the physical content of the invariant functions in the right hand side of Eq.(6). The lowest intermediate state of the correlation function $\Pi(p)$ is expressed by the baryon propagator: $\Pi(p) = -\lambda_B^2/[g^\rho - M_B - \Sigma(p, F)]$, where $M_B$ is the baryon mass, and $\lambda_B$ the coupling strength of the interpolating field to the baryon. $\Sigma(p, F)$ is the self-energy part under the influence of the background electromagnetic field, and can be decomposed into possible types of Lorentz structure. The invariant functions of the self energy on the mass-shell ($p^2 = M_B^2$) are related to physical quantities of a baryon, such as the magnetic moment, the polarizabilities, via an effective Lagrangian introduced by L’vov [1]. This Lagrangian, which was constructed by requiring Lorentz and gauge invariance, and the discrete symmetries, describes the low-energy interaction of spin-1/2 composite particles with a soft electromagnetic field. The amplitude of low energy Compton scattering is reproduced.
from the Lagrangian. For the details see our preceding paper Ref.\[12\]. We find that the invariant functions $\Pi_S$ and $\Pi_V$ are related with the polarizabilities in the following way:

$$
\Pi_S(t) = -\lambda_B^2 \left\{ \frac{1}{(t-M_B^2)^2} [4M_B^3 \Sigma_1(t)] + \frac{1}{(t-M_B^2)^2} [3M_B \Sigma_1(t) + 2M_B^2 \Sigma_S(t)] + \frac{1}{t-M_B^2} [\Sigma_S(t)] \right\},
$$
$$
\Pi_S(t) = -\lambda_B^2 \left\{ \frac{1}{(t-M_B^2)^2} [-8M_B \Sigma_1(t)] + \frac{1}{(t-M_B^2)^2} [M_B \Sigma_2(t) + M_B^3 \Sigma_3(t) + 2 \Sigma_V(t)] + \frac{1}{t-M_B^2} [M_B \Sigma_3(t)] \right\},
$$
$$
\Pi_V(t) = -\lambda_B^2 \left\{ \frac{1}{(t-M_B^2)^2} [4M_B^2 \Sigma_1(t)] + \frac{1}{(t-M_B^2)^2} [\Sigma_1(t) + 2M_B \Sigma_S(t)] \right\},
$$
$$
\Pi_V(t) = -\lambda_B^2 \left\{ \frac{1}{(t-M_B^2)^2} [4 \Sigma_1(t) - M_B^2 \Sigma_2(t)] + \frac{1}{t-M_B^2} [-\Sigma_2(t) - \frac{\Sigma_V(t)}{M_B}] \right\},
$$
$$
\Pi_V(t) = -\lambda_B^2 \left\{ \frac{1}{(t-M_B^2)^2} [-8 \Sigma_1(t)] + \frac{1}{(t-M_B^2)^2} [2 \Sigma_2(t) + M_B^2 \Sigma_3(t) + 2 \Sigma_V(t)] + \frac{1}{t-M_B^2} [\Sigma_3(t)] \right\}.
$$

$\Sigma_V$ and $\Sigma_S$ on the mass-shell are related with the polarizabilities $\bar{\alpha}_B$ and $\bar{\beta}_B$:

$$
\Sigma_V(M_B^2) = (\bar{\alpha}_B + \bar{\beta}_B)/2, \quad \Sigma_S(M_B^2) = -\bar{\beta}_B/4,
$$

while $\Sigma_1$, $\Sigma_2$ and $\Sigma_3$ with other quantities \[^{[12]}\]. Defining the following combination:

$$
\Pi^{(\bar{\alpha}B)}(p^2) = -2 \Pi_S(p^2) - p^2 \Pi_S(p^2) + 2M_B \Pi_V(p^2) + M_B \Pi_V(p^2) + p^2 M_B \Pi_V(p^2),
$$

we find the imaginary part of $\Pi^{(\bar{\alpha}B)}$ to be related solely with $\bar{\alpha}_B$:

$$
\frac{1}{\pi} \text{Im} \ \Pi^{(\bar{\alpha}B)}(t) = -\lambda_B^2 \delta(t-M_B^2) \frac{\bar{\alpha}_B}{2}.
$$

For the magnetic polarizabilities $\bar{\beta}_B$, it seems difficult to derive relations free from the constants which are not known well. Numerical estimation of $\bar{\beta}_B$ is therefore left for a future work.

We write the interpolating fields for the $n$, $\Xi^0$, $\Sigma^0$, and $\Lambda$, respectively, as follows \[^{[4]}\]:

$$
\eta_n(x) = \epsilon_{abc} [\bar{d}^a(x) C \gamma_\mu d^b(x)] \gamma_5 \gamma^\mu u^c(x),
$$
$$
\eta_{\Xi^0}(x) = -\epsilon_{abc} \left[ s^a(x) C \gamma_\mu s^b(x) \right] \gamma_5 \gamma^\mu u^c(x),
$$
$$
\eta_{\Sigma^0}(x) = \sqrt{2} \epsilon_{abc} \left\{ [u^a(x) C \gamma_\mu s^b(x)] \gamma_5 \gamma^\mu d^c(x) + [d^a(x) C \gamma_\mu s^b(x)] \gamma_5 \gamma^\mu u^c(x) \right\},
$$
$$
\eta_{\Lambda}(x) = \sqrt{2} \epsilon_{abc} \left\{ [u^a(x) C \gamma_\mu s^b(x)] \gamma_5 \gamma^\mu d^c(x) - [d^a(x) C \gamma_\mu s^b(x)] \gamma_5 \gamma^\mu u^c(x) \right\},
$$

4
where $C$ denotes the charge conjugation and $a$, $b$ and $c$ color indices. Following Ref.\cite{12} the calculations of the Wilson coefficients are straightforward.

Let us list the vacuum expectation values of interest and then classify them according to their dimensions. We are interested in the quadratic terms in $F_{\mu\nu}$. Hence, the vacuum expectation value (VEV) of the lowest-dimensional operator $\bar{q}q$ is the second order term of the quark condensate in the external field, when we expand it in powers of $F_{\mu\nu}$.

Four-dimensional operators are $F_{\mu\nu}F_{\rho\sigma}$ itself, $i\bar{q}S\gamma_\mu D_\nu q$ and $G^a_{\mu\nu}G^a_{\rho\sigma}$, where the covariant derivative is given by $D_\mu = \partial_\mu + ig(\lambda^a/2)A^a_\mu$, and $S$ is a symbol which makes the operators symmetric and traceless with respect to the Lorentz indices. The contribution of the VEV of the last two operators are so small that they can be safely neglected\cite{12}. The following three operators: $\bar{q}\sigma_{\mu\nu}q\cdot F_{\rho\sigma}$, $g\bar{q}\sigma_{\mu\nu}(\lambda^a/2)G^a_{\rho\sigma}q$, $\bar{q}S D_{\mu}D_{\nu}q$ have dimension 5. The contribution of the second operator turns out to vanish, as is for the sum rule for the masses of octet baryons. The Wilson coefficient of the last operator is constant and vanishes after a Borel transform. There are three six-dimensional operators; that is, $g\bar{q}\gamma_\mu D_\nu(\lambda^a/2)G^a_{\rho\sigma}q$, $\bar{q}q\bar{q}q$ and $\bar{q}\sigma_{\mu\nu}\bar{q}\sigma_{\rho\sigma}q$. The VEV of the first operator cannot be induced by the second-order interaction of the vacuum with an electromagnetic field, since $\bar{q}\gamma_\mu D_\nu(\lambda^a/2)G^a_{\rho\sigma}q$ has odd C-parity. To estimate the VEV’s of the last two operators we have applied the factorization hypothesis, which is firmly based in the case of the four quark operators.

$\langle \bar{q}q \rangle_F = 2\langle \bar{q}q \rangle_F \cdot \langle \bar{q}q \rangle_0,$

$\langle \bar{q}\sigma_{\mu\nu}\bar{q}\sigma_{\rho\sigma}q \rangle_F = \langle \bar{q}\sigma_{\mu\nu}q \rangle_F \cdot \langle \bar{q}\sigma_{\rho\sigma}q \rangle_F,$

where $\langle \cdot \cdot \cdot \rangle_F$ denotes the VEV in the electromagnetic field and $\langle \cdot \cdot \cdot \rangle_0$ the VEV when the external field is absent.

For the calculation up to six-dimensional terms, we are then left with the two external-field-induced VEV's: $\langle \bar{q}\sigma_{\mu\nu}q \rangle_F$, $\langle \bar{q}q \rangle_F$. It is assumed in Ref.\cite{13} that the first of them can be written, in general,

$$\langle \bar{q}\sigma_{\mu\nu}q \rangle_F = \sqrt{4\pi\alpha} F_{\mu\nu} \chi e_q \langle \bar{q}q \rangle_0,$$

where $\alpha$ is the fine structure constant and $\chi$ the so-called magnetic susceptibility of quark condensate. We extend this hypothesis to the case of the second order in $F_{\mu\nu}$ and define a “susceptibility” $\phi$ as follows:

$$\langle \bar{q}q \rangle_F = (1 + 4\pi\alpha F_{\mu\nu}F^{\mu\nu} \phi e^2_q) \langle \bar{q}q \rangle_0.$$  \hspace{1cm} (20)

In order to obtain the desired sum rule, we substitute Eq.(14) into Eq.(6). The contributions to the imaginary parts from higher mass states are approximated using the logarithmic terms in the OPE, starting at a threshold $S_0$. We calculate the left hand side of Eq.(6) by an OPE up to dimension 6. For hyperons, we further introduce the effect of the finite mass of the strange quark, $m_s$. In the sum rule for $\bar{\alpha}_{\Xi^0}$ the leading terms of the mass correction

\footnote{In a recent work\cite{15}, it is argued that $\langle \bar{q}q \rangle_F$ cannot be expanded in powers of $F_{\mu\nu}$ for massless quarks. We find, however, that $\langle \bar{q}q \rangle_F$ can be expanded in the NJL model for massless quarks\cite{16}. In this paper we assume that the expansion is possible.}
appear in the dimension 6 terms. In the sum rules for $\bar{\alpha}_\Lambda$ and $\bar{\alpha}_{\Sigma^0}$ we consider only the leading correction, which appears in the dimension 4 terms, namely $m_s \langle \bar{q}q \rangle$. One finally obtain the following results:

\[
\frac{\alpha}{4\pi^3} \left\{ \left[ 8\pi^2 \phi e_u^2 \langle \bar{u}u \rangle_0 s^2 E_1(S_0, s) + \frac{4}{3} \pi^2 \{ \chi(6e_u e_d + e_u^2) \} \langle \bar{u}u \rangle_0 sE_0(S_0, s) \right] \\
- M_\Lambda \left[ 2e_u e_d sE_0(S_0, s) - \frac{32}{3} \pi^4 \{ 4\phi e_u^2 - \chi^2 e_d^2 \} (\langle \bar{d}d \rangle_0)^2 \right] \right\} \\
= -\frac{1}{2} \lambda_n^2 \bar{\alpha}_n e^{-M_\Lambda^2/s}, 
\]

\[
\frac{\alpha}{4\pi^3} \left\{ \left[ 8\pi^2 \phi e_u^2 \langle \bar{s}s \rangle_0 s^2 E_1(S_0, s) + \frac{4}{3} \pi^2 \{ \chi(6e_u e_s + e_u^2) \} \langle \bar{s}s \rangle_0 sE_0(S_0, s) \right] \\
- M_\Sigma \left[ (e_u + e_d)e_s - 8\pi^2 \phi m_s e_s^2 \langle \bar{s}s \rangle_0 \right] sE_0(S_0, s) \\
- \frac{32}{3} \pi^4 \{ 2\phi (e_u^2 + e_d^2) - \chi^2 e_u e_d \} \langle \bar{u}u \rangle_0 (\bar{d}d)_0 \right\} \\
= -\frac{1}{2} \lambda_S^2 \bar{\alpha}_{\Sigma^0} e^{-M_\Sigma^2/s}, 
\]

\[
\frac{\alpha}{12\pi^3} \left\{ \left[ 8\pi^2 \phi (2e_u^2 \langle \bar{u}u \rangle_0 + 2e_d^2 \langle \bar{d}d \rangle_0 - e_s^2 \langle \bar{s}s \rangle_0) s^2 E_1(S_0, s) \right] \\
+ \frac{4}{3} \pi^2 \chi \left\{ 2(e_u + 3e_d + 3e_s)e_u \langle \bar{u}u \rangle_0 + 2(e_d + 3e_u + 3e_s)e_d \langle \bar{d}d \rangle_0 \\
- (e_s + 3e_u + 3e_d)e_s \langle \bar{s}s \rangle_0 \right\} sE_0(S_0, s) \right\} \\
- M_\Lambda \left[ (e_s e_u + e_s e_d + 4e_u e_d + 8\pi^2 \phi m_s (2e_u^2 \langle \bar{u}u \rangle_0 + 2e_d^2 \langle \bar{d}d \rangle_0 - 3e_s^2 \langle \bar{s}s \rangle_0) \right] \right\} sE_0(S_0, s) \\
- \frac{64}{3} \pi^4 \phi \left\{ 2(e_u^2 + e_d^2) \langle \bar{u}u \rangle_0 \langle \bar{s}s \rangle_0 + 2(e_s^2 + e_d^2) \langle \bar{d}d \rangle_0 \langle \bar{s}s \rangle_0 - (e_u^2 + e_d^2) \langle \bar{u}u \rangle_0 \langle \bar{d}d \rangle_0 \right\} \\
+ \frac{32}{3} \pi^4 \chi^2 \left\{ 2e_u e_s \langle \bar{u}u \rangle_0 \langle \bar{s}s \rangle_0 + 2e_d e_s \langle \bar{d}d \rangle_0 \langle \bar{s}s \rangle_0 - e_s e_d \langle \bar{u}u \rangle_0 \langle \bar{d}d \rangle_0 \right\} \right\} \\
= -\frac{1}{2} \lambda_{\Lambda}^2 \bar{\alpha}_\Lambda e^{-M_\Lambda^2/s}, 
\]

where $E_n(S_0, s) = 1 - \exp(-S_0/s) \sum_{k=0}^n (S_0/s)^k/k!$. It is well known that from $\Pi^{(0)}(p)$ in Eq. (2) one can get the Borel sum rules for $\lambda_n^2$ listed below.

\[
s^3 E_2(S_0, s) + \pi^2 \left( \frac{G^2}{\pi} \right) sE_0(S_0, s) + \frac{64}{3} \pi^4 (\langle \bar{d}d \rangle_0)^2 = 2(2\pi)^4 \lambda_n^2 e^{-M_\Lambda^2/s}, 
\]

\[
s^3 E_2(S_0, s) + \pi^2 \left( \frac{G^2}{\pi} \right) sE_0(S_0, s) + \frac{64}{3} \pi^4 (\langle \bar{d}d \rangle_0)^2 = 2(2\pi)^4 \lambda_S^2 e^{-M_\Sigma^2/s},
\]
For the susceptibilities we use the values evaluated in Ref.\cite{12}: 

\[ \chi = -2.58 \text{ GeV}^{-2}, \quad \phi = 1.66 \text{ GeV}^{-4}. \]  

We note, however, that there is an uncertainty due to the variation of the susceptibilities with the change of \( s \): In the reliable interval of \( s \), \( \chi \) varies from \(-2.76 \text{ GeV}^{-2}\) to \(-2.44 \text{ GeV}^{-2}\) and \( \phi \) from \(1.49 \text{ GeV}^{-4}\) to \(1.91 \text{ GeV}^{-4}\). The vacuum condensates of the operators and the mass of the strange quark are taken from Ref.\cite{14}, and are

\[ \langle \bar{u}u \rangle_0 = \langle \bar{d}d \rangle_0 = (-225 \text{ MeV})^3, \quad \langle \bar{s}s \rangle_0 = 0.8 \langle \bar{u}u \rangle_0, \quad \langle \frac{\alpha_s}{\pi} G^2 \rangle_0 = (360 \text{ MeV})^4, \quad m_s = 150 \text{ MeV}. \]

In Fig.1 we show the squared Borel mass \( s \) dependence of \( \bar{\alpha}_B \). We see that, for \( \bar{\alpha}_n \), \( \bar{\alpha}_{\Xi^0} \) and \( \bar{\alpha}_\Lambda \), the plateau develops for the suitable values of the continuum threshold parameter \( S_0 \). This implies that the sum rules for \( \bar{\alpha}_n \), \( \bar{\alpha}_{\Xi^0} \) and \( \bar{\alpha}_\Lambda \) work very well and are reliable. Indeed, we see that dominant contributions in the OPE come from the leading terms and the higher dimensional terms are suppressed. The convergence of the OPE hence appears to be good. Taking the maximum value of the curves, we predict that

\[ \bar{\alpha}_n = 14.0 \times 10^{-4} \text{fm}^3, \]  

\[ \bar{\alpha}_{\Xi^0} = 10.8 \times 10^{-4} \text{fm}^3, \]  

\[ \bar{\alpha}_\Lambda = 6.2 \times 10^{-4} \text{fm}^3. \]
These results have uncertainties arising from the errors in the susceptibilities. With the variation of the susceptibilities, \( \bar{\alpha}_n \) varies from 12.1 to 16.8, \( \bar{\alpha}_{\Xi^0} \) 9.1 to 13.5, and \( \bar{\alpha}_\Lambda \) 5.0 to 8.0.

On the other hand, the Borel curve of \( \bar{\alpha}_{\Sigma^0} \) are not stabilized for any choice of \( S_0 \) due to apparently slow convergence of the OPE. Indeed, the dimension 6 term is significantly large and contributes to \( \bar{\alpha}_{\Sigma^0} \) negatively. In view of this, no definite conclusion for \( \bar{\alpha}_{\Sigma^0} \) can be drawn. However, we can say that \( \bar{\alpha}_{\Sigma^0} \) might take a small value.

It is interesting to compare our results with those obtained in the heavy baryon chiral perturbation theory (HBChPT) \[8\]. The QCD sum rules give smaller values for the electric polarizabilities of hyperons as compared with that of neutron. This is in agreement with the HBChPT prediction. In the HBChPT, as already mentioned, the splitting is due to the contribution of kaon cloud. It can be seen that the ordering of magnitudes for strange baryons is reversal in the two calculations; namely, in the HBChPT we find \( \bar{\alpha}_{\Sigma^0} > \bar{\alpha}_{\Lambda} > \bar{\alpha}_{\Xi^0} \), while \( \bar{\alpha}_{\Xi^0} > \bar{\alpha}_{\Lambda} > \bar{\alpha}_{\Sigma^0} \) in the QCD sum rules. It is known that the effect of the decouplet states is large, so that it is highly desired to obtain the result of the HBChPT with including the decouplet states \[17\].

For charged baryons, there newly appear in the phenomenological side a triple pole term which is proportional to squared charge of the baryon and involves large factor. Within the calculation up to dimension 6 in the OPE, this term makes the Borel stability worse. We therefore need to evaluate the higher dimensional terms. This task is presently under way.

In summary, we have presented a calculation of the electric polarizabilities of neutral baryons in the framework of the QCD sum rule. The operators up to dimension 6 were retained in the operator product expansion. The sum rule predicts that the ordering of the polarizabilities is \( \bar{\alpha}_n > \bar{\alpha}_{\Xi^0} > \bar{\alpha}_{\Lambda} > \bar{\alpha}_{\Sigma^0} \), and implies that \( \bar{\alpha}_{\Sigma^0} \) has a possibility to be very small. In particular, we find that in the flavor SU(3) symmetric limit, \( \bar{\alpha}_n = \bar{\alpha}_{\Xi^0} = 4\bar{\alpha}_\Lambda/3 = 4\bar{\alpha}_{\Sigma^0} \), up to dimension 5 terms in the OPE. The calculations performed in the HBChPT have led still different predictions. However, the future experimental data from CERN and Fermilab could be of help to discriminate the different approaches.

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Fig. 1

The squared Borel mass $s$ dependence of $\bar{\alpha}_B$. The values of the continuum threshold parameter $S_0$ are taken to be $3.6 \text{ GeV}^2$ for the $n$, $4.2 \text{ GeV}^2$ for the $\Xi^0$, $6.7 \text{ GeV}^2$ for the $\Lambda$ and $\infty \text{ GeV}^2$ for the $\Sigma^0$. 
\[ \bar{\alpha}_B \left(10^{-4} \text{fm}^3\right) \] vs.

\[ s \text{ (GeV}^2\text{)} \]

Fig. 1