Gluon dipole penguin contributions to $\epsilon'/\epsilon$ and CP violation in Hyperon decays in the Standard Model

N.G. Deshpande$^1$, Xiao-Gang He$^1$ and S. Pakvasa$^2$

$^1$Institute of Theoretical Science
University of Oregon
Eugene, OR 97403-5203, USA

and

$^2$Department of Physics and Astronomy
University of Hawaii
Honolulu, Hawaii, HI 96822, USA

(Jan. 1994)

Abstract

We consider the gluon dipole penguin operator contributions to $\epsilon'/\epsilon$ and CP violation in hyperon decays. It has been proposed by Bertolini et al. that the contribution to $\epsilon'/\epsilon$ may be significant. We show that there is a cancellation in the leading order contribution and this contribution is actually suppressed by a factor of order $O(m_{\pi}^2, m_K^2)/\Lambda^2$. We find that the same operator also contributes to CP violation in hyperon decays where it is not suppressed. The gluon dipole penguin operator can enhance CP violation in hyperon decays by as much as 25%.
In this paper we study the gluon dipole penguin operator \( \bar{s}\sigma^{\mu\nu}t^a G^a_{\mu\nu}(1-\gamma_5)d \) contributions to \( \epsilon'/\epsilon \) and CP violation in hyperon decays in the Standard Model (SM). The effective \( \Delta S = 1 \) Hamiltonian at leading order can be parametrized as

\[
H_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{ud} V_{us}^* \sum_i C_i(\mu) Q_i(\mu) ,
\]

where the sum is over the effective operators \( i = 1 - 10 \) defined in Ref. [1], and the operators

\[
Q_{11} = \frac{g_s}{16\pi^2} m_s \bar{s}\sigma^{\mu\nu} t^a G^a_{\mu\nu}(1-\gamma_5)d ,
Q_{12} = \frac{eQ_d}{16\pi^2} m_s \bar{s}\sigma^{\mu\nu} F_{\mu\nu}(1-\gamma_5)d ,
\]

where \( G^a_{\mu\nu} \) and \( F^{\mu\nu} \) are the gluon and photon field strengths, respectively. \( t^a \) is the \( SU(3)_C \) generators normalized as \( \text{Tr}(t^a t^b) = \delta^{ab}/2 \). \( C_i = z_i + y_i \tau \) with \( \tau = -V_{td} V_{ts}^* / V_{ud} V_{us}^* \). CP violation is proportional to \( y_i \). The QCD corrected coefficients \( y_i \) in the SM have been evaluated in Ref. [1] and Ref. [2]. In our later calculation we will use the values in Ref. [2] for \( y_i \). The contributions to \( \epsilon'/\epsilon \) and CP violation in hyperon decays [3,4] from the operators \( Q_1 - Q_{10} \) have been extensively studied before. The contributions to \( \epsilon'/\epsilon \) from \( Q_{11,12} \) have been considered recently by Bertolini et al. [2]. The dominant contributions come from the internal top quark. It was claimed that the contribution from \( Q_{11} \) is sizable, while the contribution from \( Q_{12} \) is negligible. In this paper we reconsider the \( Q_{11} \) contribution to \( \epsilon'/\epsilon \) and also consider the contribution to CP violation in hyperon decays. We find that there is a cancellation in the contribution to \( \epsilon'/\epsilon \) and the result obtained in Ref. [2] is strongly suppressed. There are substantial contributions to CP violation in some of the hyperon decays.

**Contribution to \( \epsilon'/\epsilon \)**

The parameter \( \epsilon'/\epsilon \) is a measure of CP violation in \( K_{L,S} \to 2\pi \) decays. It is defined as

\[
\frac{\epsilon'}{\epsilon} = i \frac{e^{i(\delta_2 - \delta_0)}}{\sqrt{2}(i\xi_0 + \bar{\epsilon})} \left| \frac{A_2}{A_0} \right| (\xi_2 - \xi_0) ,
\]

where \( \bar{\epsilon} \approx 2.27 \times 10^{-3} e^{i\pi/4} \) is the CP violating parameter in \( K^0 - \bar{K}^0 \), \( \delta_i \) are the strong rescattering phases, \( |ReA_2/ReA_0| \approx 1/22 \), and \( \xi_i = ImA_i/ReA_i \). Here \( A_0 \) and \( A_2 \) are the
decay amplitudes with $I = 0$ and 2 in the final states, respectively. In order to evaluate the $Q_{11}$ contribution to $\epsilon'/\epsilon$, one needs to calculate the matrix element $< \pi^0 \pi^0 | Q_{11} | \bar{K}^0 >$. $Q_{11}$ is a $\Delta I = 1/2$ operator which contributes to $ImA_0$ only. If one uses the naive PCAC result as in Ref. [2], one obtains

$$< \pi^+ \pi^- | Q_{11} | \bar{K}^0 >= \frac{g_s m_s}{16 \pi^2 f^2 \bar{K}} \left[ < 0 | \bar{d} \sigma_{\mu\nu} t^a G^{\mu\nu}_a d | 0 > + < 0 | \bar{s} \sigma_{\mu\nu} t^a G^{\mu\nu}_a s | 0 > \right] ,$$

where $f_\pi = 132 MeV$ and $f_K = 161 MeV$. The matrix element on the right-hand side of eq.(3) can be related to the quark condensation by

$$g_s < 0 | \bar{q} \sigma_{\mu\nu} t^a G^{\mu\nu}_a q | 0 >= m_0^2 < 0 | \bar{q} q | 0 > ,$$

where $m_0^2 \approx 1 GeV^2$ is a phenomenological constant. Using

$$< 0 | \bar{u} u + \bar{s} s | 0 >= - \frac{f_K m_s^2}{m_u + m_s} ,$$

and $< 0 | \bar{d} d | 0 >= < 0 | \bar{u} u | 0 >$, one finally obtains

$$< \pi^+ \pi^- | Q_{11} | \bar{K}^0 >= - \frac{m_0^2 m_s}{16 \pi^2 (m_u + m_s)} \frac{m_K^2 f_K}{f^2} .$$

One also has, $< \pi^0 \pi^0 | Q_{11} | \bar{K}^0 >= < \pi^+ \pi^- | Q_{11} | \bar{K}^0 >$. Using the matrix element in eq.(7), a quite large contribution to $\epsilon'/\epsilon$ was obtained in Ref. [2]. They find that $Q_{11}$ contributes between $(2 \sim 3) \times 10^{-4}$. It weakly depends on the top quark mass $m_t$ for $m_t$ between 100 GeV to 250 GeV. We would like to point out that the naive PCAC result obtained above is not correct. The calculation in Ref. [2] is only part (Fig. 1.a) of the contributions. An important "tadpole" contribution (Fig. 1.b) was not considered in the analysis of Ref. [2]. This contribution cancels exactly the PCAC result obtained above. The net contribution is much smaller. The situation is the same as for $\epsilon'/\epsilon$ in the Weinberg model of CP violation [7,8]. The importance of the "tadpole" contribution was first noticed by Donoghue and Holstein [7] in the Weinberg model of CP violation. We now calculate these contributions in the SM.

The leading order chiral realization of $Q_{11}$ for $\bar{K}^0 \to n\pi$ is
\[ L = a Tr(\lambda_6 + i\lambda_7)U) + H.C. \quad (8) \]

The conventional parametrization of \( U \) is given by \( U = \exp(-i4t^a\phi^a/f_\pi) \) with \( \phi^a \) being the fields of the pseudoscalars. In our discussion we follow Ref. \[8\] using a more general parametrization of \( U \). To fourth power in the meson fields, there is a free parameter \( a_3 \) \[8,9\],

\[
U = 1 + i \frac{2}{f_\pi}(2t^a\phi^a) - \frac{2}{f_\pi^2}(2t^a\phi^a)^2 - i \frac{a_3}{f_\pi^3}(2t^a\phi^a)^3 + \frac{2(a_3 - 1)}{f_\pi^4}(2t^a\phi^a)^4 + ... (9)
\]

The conventional parametrization corresponds to \( a_3 = 4/3 \). The on-shell amplitudes should not depend on \( a_3 \). The effective lagrangian in eq.(8) will not only generate the direct \( K \to 2\pi \) amplitude (Fig. 1.a), but also generate a non-zero \( K \to \text{vacuum} \) transition amplitude. This non-zero \( K \to \text{vacuum} \) transition amplitude, when combined with the strong \( K\pi\bar{K}\pi \) amplitude \( A_{\text{strong}}(K\bar{K}\pi^0\pi^0) \) with one \( K \) off-shell with \( q^2 = 0 \), will also generate a \( K \to 2\pi \) amplitude as show in Fig. 1.b. The total amplitude for \( K \to 2\pi \) from \( Q_{11} \) is the sum of contributions from Fig. 1.a and Fig. 1.b. We have

\[
A_{\text{total}}(K^0 \to \pi^0\pi^0) = A(\text{fig.1.a}) + A(\text{fig.1.b})
\]

\[
= A(K^0 \to \pi^0\pi^0)|_{\text{fig.1.a}} + A_{\text{strong}}(K^0\bar{K}^0\pi^0\pi^0) \frac{1}{m_K^2} A(K^0 \to \text{vacuum}) , (10)
\]

where \( A(K^0 \to n\pi) = < n\pi | - H_{\text{eff}}(Q_{11}) | K^0 > \), and \( H_{\text{eff}}(Q_{11}) = -i \frac{G_F}{\sqrt{2}} y_{11} Im(V_{td}V_{ts}^*) Q_{11} \). We have

\[
A(K^0 \to \text{vacuum}) = -i \frac{G_F}{\sqrt{2}} y_{11} Im(V_{td}V_{ts}^*) f_\pi \sqrt{\frac{g_s m_s}{32\pi^2}} \tilde{A}_{K\pi} ,
\]

\[
A(K^0 \to \pi^0) = -i \frac{G_F}{\sqrt{2}} y_{11} Im(V_{td}V_{ts}^*) \frac{g_s m_s}{32\pi^2} \tilde{A}_{K\pi} ,
\]

\[
A(K^0 \to \pi^0\pi^0) = i \frac{G_F}{\sqrt{2}} y_{11} Im(V_{td}V_{ts}^*) \frac{g_s m_s}{32\sqrt{2}\pi^2 f_\pi} \tilde{A}_{K\pi} \frac{a_3}{2} . (11)
\]

Here \( \tilde{A}_{K\pi} = -i < \pi^0 s\gamma^\mu 2t^a G_{\mu\nu} (1 - \gamma_5)d | K^0 > \). Note that \( A(K^0 \to \pi^0\pi^0) \) is \( a_3 \) dependent, which can not be the final answer. Additional contribution from Fig. 1.b has to be considered. What has been evaluated in Ref. \[2\] corresonds to \( A(\text{fig.1.a}) \) with \( a_3 = 2 \). \( m_s \tilde{A}_{K\pi} \) has been calculated to be \( 0.11 \sim 0.17 \ GeV^2 \) in the MIT bag model \[10\]. Using this value one obtains approximately the same numerical value for \( A(\text{fig.a}) \) as obtained in Ref. \[2\].

\[4\]
To calculate the contribution from Fig. 1.b, one needs to know the strong $K\pi\bar{K}\pi$ amplitude $A_{\text{strong}}$. This can be obtained from the leading chiral lagrangian

\[
L = \frac{f_\pi^2}{8} [ Tr \partial_\mu U \partial^\mu U + B Tr (MU + U^\dagger M) ] ,
\]

(12)

where $M = \text{Diag}(m_u, m_d, m_s)$ is the quark mass matrix and $B$ is a constant. From this we obtain [8,9]

\[
A_{\text{strong}}(K^0\bar{K}^0\pi^0\pi^0) = \frac{m_K^2}{2f_\pi^2} \frac{a_3}{2} .
\]

(13)

Here we have set the momentum of $K$ annihilated into the vacuum to be zero and others to be on-shell (Of course the on-shell $A_{\text{strong}}$ is $a_3$ independent.). Inserting eq.(13) into eq.(10), we find the $a_3$ dependent contribution from Fig. 1.b cancels exactly that from Fig. 1.a. To this order, there is no contribution to $\epsilon'/\epsilon$ from $Q_{11}$. However, the cancellation may not be complete due to higher derivative terms in chiral perturbation theory. Unfortunately, one does not know how to calculate higher-order contributions to the matrix elements. One can define a suppression factor $D$,

\[
A_{\text{total}}(K^0 \rightarrow \pi^0\pi^0) = A(\text{fig.1.a)}D .
\]

(14)

The suppression factor should be of order $O(p^2)/\Lambda^2 \approx O(m_K^2, m_\pi^2)/\Lambda^2 [8]$. Here $\Lambda \approx 1 \text{ GeV}$ is the chiral symmetry breaking scale. The contribution is suppressed by a factor of 0.3 or even more. The contribution from $Q_{11}$ calculated in Ref. [2] correspond to $D = 1$.

We list the results for $D = m_K^2/\Lambda^2$ in Table 1, but note that the sign of the contribution is unknown. The parameter $\text{Im}(V_{td}V_{ts}^*)$ is constrained by the parameter $\epsilon$ from $K^0 - \bar{K}^0$ mixing. Unfortunately it is not completely fixed. It can vary between $3 \times 10^{-4}$ to $10^{-4}$ for top quark mass $m_t$ about 100 GeV and $2 \times 10^{-4}$ to $0.5 \times 10^{-4}$ for $m_t$ about 200 GeV [11]. In Table 1, following Ref. [2], we use an approximation $\text{Im}(V_{td}V_{ts}^*) = 2.77 \times 10^{-4}(m_t^2/m_W^2)^{-0.365}$ as the central value. We see that the effect of $Q_{11}$ is generally insignificant, but becomes important when the contribution from $Q_1 - Q_{10}$ become small due to cancellations for $m_t$ around 200 GeV.
CP violation in hyperon decays

CP violation in hyperon decays in the SM has been studied before [3–5]. The $Q_{11}$ contributions were not included, and we now turn to study these contributions. Non-leptonic hyperon decays proceed into both S-wave (parity-violating) and P-wave (parity-conserving) final states with amplitudes $S$ and $P$, respectively. One can write the amplitude in the rest frame of the initial baryon as

$$Amp(B_i \to B_f \pi) = S + P \vec{\sigma} \cdot \vec{q},$$

where $\vec{q}$ is the momentum of pion. It is convenient to write the amplitudes as

$$S = \sum_i S_i e^{i(\phi_i^S + \delta_i^S)}$$
$$P = \sum_i P_i e^{i(\phi_i^P + \delta_i^P)}$$

(16)

to explicitly separate the strong rescattering phases $\delta_i$ and the weak CP violating phases $\phi_i$. In the rest frame of the initial baryon, one particularly interesting observable is the asymmetry $A$ defined in Ref. [4]

$$A = \frac{\alpha + \bar{\alpha}}{\alpha - \bar{\alpha}},$$

(17)

where $\alpha = 2Re(S^*P)/(|S|^2 + |P|^2)$, and $\bar{\alpha}$ is the corresponding quantity for anti-hyperon decays. A non-zero $A$ signals CP violation. In this paper we will concentrate on the study of $A$. The results can be easily generalized to other CP violating observables defined in Ref. [4]. We will calculate the CP violating observable $A$ for $\Lambda \to N\pi$ and $\Xi \to \Lambda\pi$. Since to leading order the observables $A(\Lambda^0)$ and $A(\Xi^-)$ for $\Lambda \to p\pi^-$ and $\Xi^- \to \Lambda\pi^-$ are the same as $A(\Lambda^0)$ and $A(\Xi^0)$ for $\Lambda \to n\pi^0$ and $\Xi^0 \to \Lambda\pi^0$, respectively, we choose to work with $\Lambda \to n\pi^0$ and $\Xi^0 \to \Lambda\pi^0$ decays.

For the same reason as for the $K \to 2\pi$ amplitude, when evaluating the S-wave amplitude, one should also include the contributions from the direct contribution of Fig. 2.a and the "tadpole" contribution of Fig. 2.b. We have [4][12]
\[ S(\Lambda \to n\pi^0) = -\frac{i}{\sqrt{2}f_\pi} < n|H_{eff}^\dagger(Q_{11})|\Lambda >_{fig.2.a} \]

\[ + \frac{i}{\sqrt{2}f_\pi} \left[ \frac{3}{2} \right]^{1/2} \frac{M_\Lambda - M_n}{m_s - m_d} \left[ \frac{m_s + m_d}{i f_K M_K^2} \sqrt{2} < 0|H_{eff}^\dagger(Q_{11})|K^0 > \right]_{fig.2.b}, \]

(18)

\[ S(\Xi^0 \to \Lambda\pi^0) = -\frac{i}{\sqrt{2}f_\pi} < \Lambda|H_{eff}^\dagger(Q_{11})|\Xi^0 >_{fig.2.a} \]

\[ - \frac{i}{\sqrt{2}f_\pi} \left[ \frac{3}{2} \right]^{1/2} \frac{M_\Xi - M_\Lambda}{m_s - m_d} \left[ \frac{m_s + m_d}{i f_K M_K^2} \sqrt{2} < 0|H_{eff}^\dagger(Q_{11})|K^0 > \right]_{fig.2.b}. \]

(19)

We use pole model to calculate the P-wave amplitudes. For consistency, one should include both the baryon and Kaon poles [12]. The "tadpole" contribution from Fig. 2.b in the S-wave amplitude and the Kaon pole contribution in the P-wave amplitude are both sizable. However, the CP violating observable A depends on the difference between the S-wave phase \( \phi^S \) and the P-wave phase \( \phi^P \). There is substantial cancellation between the "tadpole" contribution in the S-wave and the Kaon pole contribution in the P-wave.

The contributions to CP violation in hyperon decays from the same operator \( Q_{11} \) have been studied in the Weinberg model [3,4]. We can use some of the results from there. However in the Weinberg model the CP violating parameter \( \bar{\epsilon} \) in the \( K^0 - \bar{K}^0 \) is also generated by the operator \( Q_{11} \) through long distance \( \pi, \eta \) and \( \eta' \) pole contributions, the coefficient \( C_{11}(Weinberg) \) of \( Q_{11} \) [7] is fixed. In our case, the contribution to \( \bar{\epsilon} \) from \( Q_{11} \) is small. To obtain the CP violating phases in the decay amplitudes, we can repeat the MIT bag model calculations in Ref. [4]. In fact we can obtain the phases by replacing the matrix element

\[ < \pi^0| - \frac{G_F}{\sqrt{2}} Im(V_{ud}^* V_{us} C_{11}^*(Weinberg)) Q_{11}^\dagger|K^0 > \]

in the Weinberg model by the SM matrix element, \( < \pi^0| - H_{eff}^\dagger(Q_{11})|K^0 > \), and rescale the phases. We define the rescaling factor \( R \) as

\[ R = \frac{< \pi^0| - H_{eff}^\dagger(Q_{11})|K^0 >}{< \pi^0| - \frac{G_F}{\sqrt{2}} Im(V_{ud}^* V_{us} C_{11}^*(Weinberg)) Q_{11}^\dagger|K^0 >} = \frac{-G_F}{\sqrt{2}} y_{11} Im(V_{td}^* V_{ts}^*) \frac{g_s A_{K\pi}}{32\pi^2} \frac{m_s \tilde{A}_{K\pi}}{5.8 \times 10^{-11} \text{ GeV}^2}. \]

(20)

Here we have used \( < \pi^0| - \frac{G_F}{\sqrt{2}} Im(V_{ud}^* V_{us} C_{11}^*(Weinberg)) Q_{11}^\dagger|K^0 > = 5.8 \times 10^{-11} \text{ GeV}^2 \) [4]. We will follow Weinberg to use \( g_s \approx 4\pi/\sqrt{6} \) [13]. The CP violating phases in hyperon decays
from $Q_{11}$ in the SM are now obtained by multiplying $R$ to the phases obtained in Ref. [3] for the Weinberg model.

In Table 2 we list the CP violating observables $A$ for $\Lambda$ and $\Xi$ decays from different operators. From Table 2, we see that the $Q_{11}$ contribution to CP violation in $\Lambda \rightarrow N\pi$ is negligible. However the contribution to CP violation in $\Xi \rightarrow \Lambda\pi$ can be important. The result is to enhance CP violation in $\Xi$ decays by about 17%. If we use the upper value of $0.17 \, GeV^4$ for $m_s \tilde{A}_{K\pi}$, the enhancement will be 25%. The allowed regions for $A(\Lambda)$ and $A(\Xi)$ without the contributions from $Q_{11}$ are $3 \times 10^{-5} \sim 0.4 \times 10^{-5}$ and $5.3 \times 10^{-5} \sim 0.7 \times 10^{-5}$. With the new contributions from $Q_{11}$ the allowed region for $A(\Lambda)$ is not changed, but the allowed region for $A(\Xi)$ becomes to $6.6 \times 10^{-5} \sim 0.7 \times 10^{-5}$. This may be tested in the future [14].

ACKNOWLEDGMENTS

This work was supported in part by the Department of Energy Grant No. DE-FG06-85ER40224 and DE-AN03-76SF00235. We thank E. Golowich for reading the manuscript.
REFERENCES

[1] J. Flynn and L. Randall, Phys. Lett. B224, 221(1989); G. Buchalla, A. Buras and Harlander, Nucl. Phys. B337, 313(1990); E.A. Paschos and Y.-L. Wu, Mod. Phys. Lett.A6, 93(1991); M. Lusignoli et al., Nucl. Phys. B369, 139(1992); A. Buras and M. Lautenbacher, Preprint, MPI-PI/93-60, TUM-T31-47/93.

[2] S. Bertolini, M. Fabbrichesi and E. Gabrielli, Reprint, CERN-TH. 7097/93; SISSA 174/93/EP.

[3] J. Donoghue and S. Pakvasa, Phys. Rev. Lett. 55, 162(1985).

[4] J. Donoghue, Xiao-Gang He and S. Pakvasa, Phys. Rev. D34, 833(1986).

[5] Xiao-Gang He, H. Steger and G. Valencia, Phys. Lett. B272, 411(1991).

[6] B. Ioffe, Nucl. Phys. B188, 317(1981); V. Belyaev and B. Ioffe, Sov. Phys. JETP 56, 493(1982); S. Nikolaev and A. Radyushkin, Nucl. Phys. B213, 285(1983); Y. Cheng, H. Dosch, M. Kremer and D. Schall, Z. Phys. C25, 151(1984); M. Kremer and G. Schierholz, Phys. Lett. B194, 283(1987); S. Narison, Phys. Lett. B210, 238(1988).

[7] J. Donoghue and B. Holstein, Phys. Rev. D32, 1152(1985).

[8] H. - Y. Cheng, Phys. Rev. 34, 1397(1986).

[9] J. Cronin, Phys. Rev. 161, 1483(1967).

[10] J. Donoghue, J. Hagelin and B. Holstein, Phys. Rev. D25, 195(1982).

[11] G. Harris and J. Rosner, Phys. Rev. D45, 946(1992); Y.-L. Wu, Int. J. Mod. Phys.A7, 2863(1992).

[12] J. Donoghue and B. Holstein, Phys. Rev. D33, 2717(1986).

[13] S. Weinberg, Phys. Rev. Lett. 63, 2333(1990).

[14] G. Gidal, P.M. Ho and K.B. Luk, Preprint, Fermilab Prposal P-871.
TABLES

TABLE I. $\epsilon'/\epsilon \times 10^4$ for $\text{Im}(V_{td}V_{ts}^*) = 2.77 \times 10^{-4}(m_t^2/m_W^2)^{-0.365}$.

| $m_t$(GeV) | 130 | 170 | 200 | 230 |
|------------|-----|-----|-----|-----|
| $\Lambda_4 = 200 \text{ MeV}$ |     |     |     |     |
| $Q_1 - Q_{10}$ | 6.3 | 2.8 | 0.7 | -1.4 |
| $Q_{11}$ | 0.65 | 0.58 | 0.5 | 0.48 |
| $\Lambda_4 = 300 \text{ MeV}$ |     |     |     |     |
| $Q_1 - Q_{10}$ | 7.8 | 3.6 | 1.0 | -1.7 |
| $Q_{11}$ | 0.73 | 0.63 | 0.55 | 0.50 |

TABLE II. $A \times 10^5$ for $\text{Im}(V_{td}V_{ts}^*) = 2.77 \times 10^{-4}(m_t^2/m_W^2)^{-0.365}$ and $m_s\bar{A}_{K\pi} = 0.12 \text{ GeV}^4$.

| $m_t$(GeV) | 130 | 170 | 200 | 230 |
|------------|-----|-----|-----|-----|
| $\Lambda_4 = 200 \text{ MeV}$ |     |     |     |     |
| $A(\Lambda_0^0)(Q_1 - Q_{10})$ | -1.5 | -1.3 | -1.1 | -1.0 |
| $A(\Lambda_0^0)(Q_{11})$ | -0.043 | -0.037 | -0.033 | -0.031 |
| $A(\Xi^-)(Q_1 - Q_{10})$ | -2.6 | -2.2 | -1.9 | -1.8 |
| $A(\Xi^-)(Q_{11})$ | -0.55 | -0.47 | -0.41 | -0.40 |
| $\Lambda_4 = 300 \text{ MeV}$ |     |     |     |     |
| $A(\Lambda_0^0)(Q_1 - Q_{10})$ | -2.0 | -1.7 | -1.5 | -1.4 |
| $A(\Lambda_0^0)(Q_{11})$ | -0.046 | -0.039 | -0.034 | -0.032 |
| $A(\Xi^-)(Q_1 - Q_{10})$ | -3.4 | -2.9 | -2.5 | -2.3 |
| $A(\Xi^-)(Q_{11})$ | -0.58 | -0.50 | -0.44 | -0.40 |
Fig. 1. Contributions to $K \to \pi\pi$ amplitude. Here $W$ and $St$ indicate weak and strong interactions, respectively.

Fig. 2. Contributions to the S-wave $B_i \to B_f \pi$ amplitudes.