SU(3) HIGGS BOSONS AND SOLITON CONTRIBUTIONS TO THE NON-LEPTONIC WEAK INTERACTIONS

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ABSTRACT

A flavor $SU(3)_f$ quark and octet Higgs boson model in one spatial dimension is studied in order to explore some dynamical processes of non-leptonic weak interactions. In this model, the d-s quark mixing and the soliton state contributions to non-leptonic weak decays are presented.

1. INTRODUCTION

The flavor mixing in the weak currents is essential for the non-leptonic weak decays. However, it seems to us that a mechanism of d-s mixing in the weak quak-currents is not yet clear. We present one possibility of the flavor mixing of quark in a flavor $SU(3)_f$ quark and octet Higgs boson model in one spatial dimension which is a variant of Jackiw-Rebbi's one[1]. It will be shown that a spontaneous broken $SU(3)_f$ symmetry of boson ground state leads to a flavor mixing of quark and also to the existence of soliton and quark-soliton states. The soliton looks very much like a classical extended particle. Its energy density is localized at a point, its total energy is finite and it is stable.

The quark-soliton system may look like an electron-nuclear Coulomb potential system. It is known that a quark in soliton potential (in one spatial dimension) has only one bound state with zero energy (quark-soliton). The quark-soliton interactions give rise to additional proccsses of thc non-leptonic weak decays. If we could extend our model to a realistic three spatial model, we may expect that these processes contribute to the enhancement of right-handed quark. For caution's sake, it should be noted that in one spatial dimension, a quark is handless.

In Sec. 2, a broken $SU(3)_f$ symmetry ground state is defined. In Sec. 3, we present the quark solutions and their quantization. Sec. 4 is devoted to estimate the soliton state contribution to the non-leptonic weak process in which the enhancement of rigthanded quark is suggested.

2. BROKEN SYMMETRY GROUND STATE

Let us consider a quantum field theory of flavor $SU(3)_f$ quark $q\equiv(u, d, s)$ and neutral, scalar octet-bosons $\phi_\alpha(\alpha=1, 2, \cdots\cdots, 8)$ in one spatial dimension. The Lagrangian density is assumed to be of the form

$$\mathcal{L}=\frac{1}{2}\partial_\mu\phi^\alpha\partial^\mu\phi_\alpha-\frac{\lambda^2}{2\mu^2}\left[\mu^2-\phi^\alpha\phi_\alpha\right]^2+i\bar{q}\gamma^\mu\partial_\mu q-g\bar{q}\lambda_\alpha q\phi^\alpha, \tag{1}$$

where $\bar{q}=q^\dagger\gamma_0$, $\gamma^0=\begin{bmatrix}0&1\\1&0\end{bmatrix}$, $\gamma^1=\begin{bmatrix}i&0\\0&-i\end{bmatrix}$ and λ_α 's are the Gell-Mann matrices.

This is a variant of Jackiw-Rebbi's model which gives the existence of zero-energy quark eigenstate, localized in the vicinity of one soliton of the Bose field.[2] In order to review the soliton states in this model, let us for the moment ignore the quark field in (1).

The field equation satisfied by ϕ_α is

$$\Box\phi_\alpha+\frac{\partial U}{\partial\phi^\alpha}=\left[\frac{\partial^2}{\partial t^2}-\frac{\partial^2}{\partial x^2}\right]\phi_\alpha+\frac{\partial U}{\partial\phi^\alpha}=0, \tag{2}$$

where

$$U(\phi_\alpha\phi^\alpha)=\frac{\lambda^2}{\mu^2}\left[\mu^2-\phi_\alpha\phi^\alpha\right]^2. \tag{3}$$

For a static solution $\phi_\alpha(x)$, the equation (2) is reduced to

$$\frac{d^2}{dx^2}\phi_\alpha+\frac{2\lambda^2}{\mu^2}\left[\mu^2-\phi_\beta\phi^\beta\right]\phi_\alpha=0. \tag{4}$$

The first integration of (4) is

$$\frac{1}{2}\frac{d\phi_\alpha}{dx}\frac{d\phi^\alpha}{dx}=U(\phi^\alpha\phi_\alpha), \tag{5}$$

and the static field energy E_s is given by

$$\begin{aligned}E_s&=\int dx\left\{\frac{1}{2}\frac{d\phi_\alpha}{dx}\frac{d\phi^\alpha}{dx}+U(\phi_\alpha\phi^\alpha)\right\}\\&=2\int dx\ U(\phi_\alpha\phi^\alpha)\geq 0.\end{aligned} \tag{6}$$

The degenerate vacua are found at the points

$$\phi_\alpha \phi^\alpha = \mu^2 , \tag{7}$$

where $U(\phi_\alpha \phi^\alpha)=0$ (see Fig. 1).

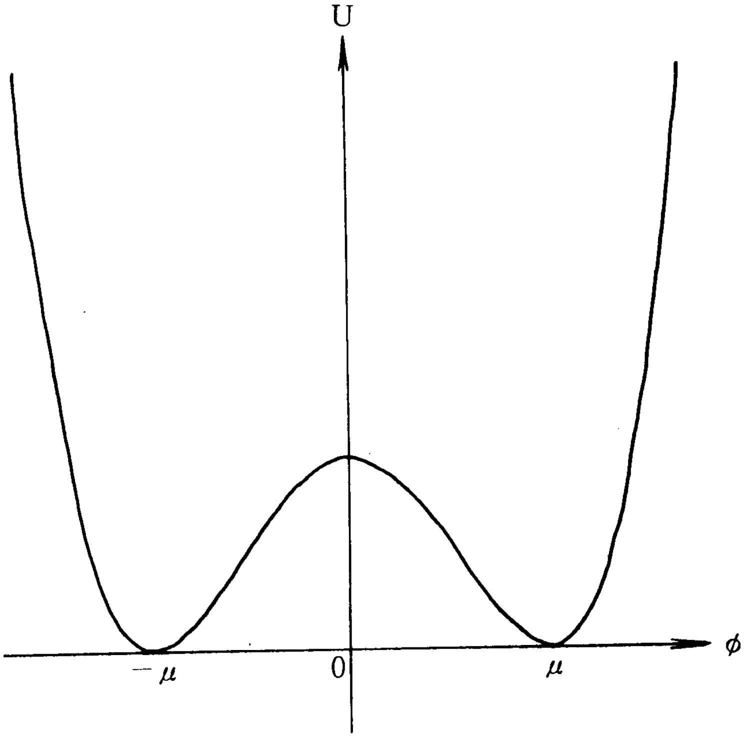

Fig. 1: $\phi = \pm\sqrt{\phi_\alpha \phi^\alpha}$
The octet boson vacuum is defined at the points $\phi = \pm\mu$.

For spontaneous breaking of $SU(3)_f$ symmetry, we assume that all vacua satisfy the conditions

$$\phi_1,\ \phi_2,\ \phi_3,\ \phi_4,\ \phi_5 = 0,$$

and

$$\phi_6 = \pm\mu \sin\theta \cos\varphi,\quad \phi_7 = \pm\mu \sin\theta \sin\varphi,$$
$$\phi_8 = \pm\mu \cos\theta. \tag{8}$$

We make choice of the plus-sign solutions for ϕ_6, ϕ_7, ϕ_8 in (8), and define our vacuum by the conditions

$$\begin{aligned} &\phi_\alpha = 0 \quad (\alpha = 1,\ 2,\ \cdots\cdots,\ 5), \\ &\phi_6 = \mu \sin\theta \cos\varphi, \\ &\phi_7 = \mu \sin\theta \sin\varphi, \\ &\phi_8 = \mu \cos\theta. \end{aligned} \tag{9}$$

We note that our vacuum is in a polarized flavor state.

We find the soliton shape of ϕ_α which is a x-dependent solution,

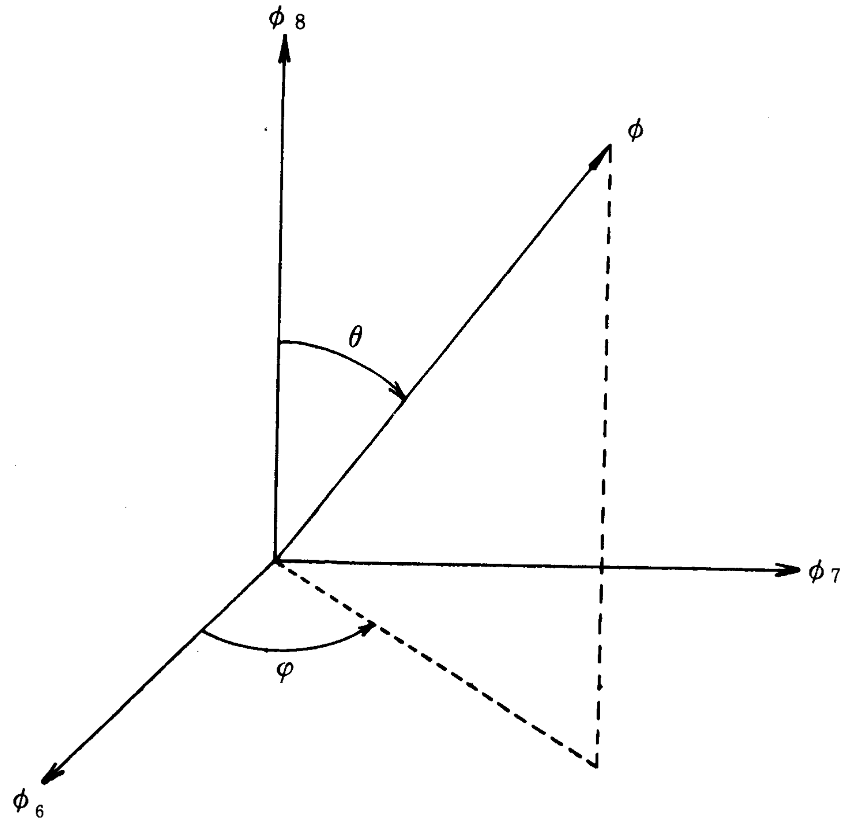

Fig. 2. The flavor polarization of vacuum.

$$\begin{aligned} &\phi_\alpha=0 \quad (\alpha=1,\ 2,\ \cdots\cdots,\ 5), \\ &\phi_6=\pm\mu\sin\theta\cos\varphi\tanh\lambda(x-x_0), \\ &\phi_7=\pm\mu\sin\theta\sin\varphi\tanh\lambda(x-x_0), \\ &\phi_8=\pm\mu\cos\theta\tanh\lambda(x-x_0) \quad , \end{aligned} \tag{10}$$

in which x_0 is the point where the soliton energy density is localized, and a soliton is translational invariant. Corresponding to our vacuum definition (9), we take plus-sign solutions.

The soliton energy is finite;

$$\begin{aligned} E_s &= \int dx\left\{\frac{1}{2}\frac{d\phi_\alpha}{dx}\frac{d\phi^\alpha}{dx}+U(\phi_\alpha\phi^\alpha)\right\} \\ &= \int dx\left[\frac{d\phi_\alpha}{dx}\frac{d\phi^\alpha}{dx}\right] = \frac{4}{3}\mu^2\lambda . \end{aligned} \tag{11}$$

3. QUARK-SOLITONS

In the vacuum sector, the quark-equation is given by

$$\{\alpha p+\beta g\mu(\lambda_6\sin\theta\cos\varphi+\lambda_7\sin\theta\sin\varphi+\lambda_8\cos\theta)\}q^{(\pm)}=\pm|\varepsilon|q^{(\pm)}, \tag{12}$$

where $\alpha=\gamma^0\gamma^1=\begin{bmatrix}0 & -i\\ i & 0\end{bmatrix}$ and $\beta=\gamma^0=\begin{bmatrix}0 & 1\\ 1 & 0\end{bmatrix}$.

In order to diagonalize the mass matrix:

$$g\mu(\lambda_6 \sin\theta \cos\varphi + \lambda_7 \sin\theta \sin\varphi + \lambda_8 \cos\theta), \tag{13}$$

we define the unitary matrix U by

$$U = \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{i\varphi}\cos\theta_c & \sin\theta_c \\ 0 & -\sin\theta_c & e^{-i\varphi}\cos\theta_c \end{pmatrix}, \tag{14}$$

with

$$\tan\theta_c = \sqrt{\frac{1-K}{1+K}}$$

and $$K = \left[1+\frac{4}{3}\tan^2\theta\right]^{-1/2}.$$

By this unitary transformation, the mass matrix (13) is diagonalized.

$$\delta m = U \cdot g\mu(\lambda_6 \sin\theta \cos\varphi + \lambda_7 \sin\theta \sin\varphi + \lambda_8 \operatorname{con}\theta) \cdot U^\dagger = \mu\tilde{g}, \tag{15}$$

where $\tilde{g} = \begin{pmatrix} g_1 & 0 & 0 \\ 0 & g_2 & 0 \\ 0 & 0 & g_3 \end{pmatrix}$,

$$g_1 = g\frac{2\cos 2\theta_c}{\sqrt{3(3+\cos 2\theta_c)}},\ g_2 = g\frac{3-\cos 2\theta_c}{\sqrt{3(3+\cos 2\theta_c)}} \text{ and } g_3 = -g\frac{3+\cos 3\theta_c}{\sqrt{3(3+\cos 2\theta_c)}}.$$

The equation (12) is reduced to

$$(\alpha p + \beta\mu\tilde{g})q'^{(\pm)} = \pm|\varepsilon|q'^{(\pm)}. \tag{16}$$

The transformed quark fields $q' = Uq$ are given by

$$\begin{aligned} q'_1 &= q_1, \\ q'_2 &= \cos\theta_c e^{i\varphi} q_2 + \sin\theta_c q_3, \\ q'_3 &= -\sin\theta_c q_2 + \cos\theta_c e^{-i\varphi} q_3, \end{aligned} \tag{17}$$

where $q_1 = u$, $q_2 = d$ and $q_3 = s$.

This shows that our spontaneous broken-symmetry vacuum leads to the d-s mixing with the Cabbibo angle θ_c.

$$\theta_c = \tan^{-1}\sqrt{\frac{1-K}{1+K}}.$$

In the soliton sector, the same unitary transformation (14) reduces the quark equation to the form

$$\{\alpha p + \beta\tilde{g}\mu \cdot \tanh\lambda(x-x_0)\}q'^{(\pm)} = \pm|\varepsilon|q'^{(\pm)} \tag{18}$$

which is Dirac's equations in the presence of the static potential

$$V(x-x_0) = \tilde{g}\mu \cdot \tanh\lambda(x-x_0).$$

The equations admit the zero-energy solutions

$$q'_{(0)i}(x-x_0) = C_i \begin{bmatrix} 1 \\ 0 \end{bmatrix} \{\cosh\lambda(x-x_0)\}^{-g_i\mu/\lambda} \quad (i=1,\ 2,\ 3). \tag{19}$$

The normalization constant C_i is given by

$$C_i = \frac{1}{\sqrt{\int_{-\infty}^{\infty} dx\{\cosh \lambda(x-x_0)\}^{-2g_i\lambda/\mu}}}\ .$$

We note that these zero-energy solutions are charge conjugation self-conjugate

$$q'_{(0)}{}^c = \sigma_3 q'_{(0)}{}^* = q'_{(0)}, \tag{20}$$

where $\sigma_3 = \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix}$,

and also the soliton-quark is localized around the point x_0 which is assumed to be the center of soliton configuration. So far we ignored other Higgs mechanisms which generate the main part of quark masses m_0. Taking account of such mass generations, the equation of (15) and (18) should be replaced by

$$\{\alpha p + \beta(m_0 + \tilde{g}\mu)\} q'^{(\pm)} = \pm |\varepsilon| q'^{(\pm)}, \tag{21}$$

and

$$\{\alpha p + \beta(m_0 + \tilde{g}\mu \cdot \tanh\lambda(x-x_0))\} q'^{(\pm)} = \pm |\varepsilon| q'^{(\pm)} \tag{22}$$

Corresponding to (19), the equation (22) have the zero-energy solutions:

$$q'_{(0)i}(x-x_0) = C'_i \begin{bmatrix} 1 \\ 0 \end{bmatrix} e^{-m_0|x-x_0|}\{\cosh\lambda(x-x_0)\}^{-g_i\mu/\lambda}, \tag{23}$$

the normalization constants being

$$C'_i = \frac{1}{\sqrt{\int_{-\infty}^{\infty} dx\, e^{-2m_0|x-x_0|}\{\cosh\lambda(x-x_0)\}^{-2g_i\mu/\lambda}}}\ . \tag{24}$$

Following Jackiw-Rebbi,[1,2] we assume that the quatization of quark fields is achieved by the usual expansion in modes. In the vacuum sector, the expansion of a quark operator $\hat{q}'_{(i)}$ is defined by

$$\hat{q}'_{(i)} = \sum_k \{e^{-i\varepsilon_k^{(i)}t} b_k^{(i)}\ u_{(i)k}^{(+)}(x) + e^{i\varepsilon_k^{(i)}t} d_k^{(i)\dagger} v_{(i)k}^{(-)}(x)\}. \tag{25}$$

$u_{(i)k}^{(+)}$ is a positive energy solution of (21), and $v_{(i)k}^{(-)}$ is the charge conjugate of the negative energy solutions. The energy spectrum is

$$|\varepsilon_k^{(i)}| = \sqrt{m_i^2 + k^2},\quad m_i = m_0 + \mu g_i\,. \tag{26}$$

The quark creation (annihilation) operators $b_k^{(i)\dagger}$ ($b_k^{(i)}$) and the anti-quark creation (annihilation) operators $d_k^{(i)\dagger}$ ($d_k^{(i)}$) obey the usual commutation relations

$$\begin{aligned} [b_k^{(i)}\ ,\ b_{k'}^{(j)\dagger}]_+ &= \delta_{ij}\,\delta(k-k'), \\ [d_k^{(i)}\ ,\ d_{k'}^{(j)\dagger}]_+ &= \delta_{ij}\,\delta(k-k'), \end{aligned} \tag{27}$$

any other anti-commutator$=0$.

In the soliton sector, we may imagine a quark to be just as a quarkic atom like a μ-mesic atom whose center is situated at x_0. Let us denote the quark field operator in soliton sectors by $\hat{q}_i'(x-x_0)$. Generally, the parameter x_0 in $\hat{q}_i'(x-x_0)$ should be a quantum variable $\hat{x}_0$ which corresponds to a position operator of collective motion of an extended particle-like soliton. However, we ignore the quantum nature of $\hat{x}_0$ and replace it by an expectation value $x_0=\langle\hat{x}_0\rangle$ in the following. Moreover, we assume that the $\hat{q}_i'(x-x_0)$'s with different value of x_0 are each other independent field operators which are specified by x_0.

Quantization again is defined by an expansion in modes

$$\hat{q}_i'(x-x_0)=a^{(i)}(x_0)q_{i\,(0)}(x-x_0)+\sum_k\{e^{-i\varepsilon_k^{(i)}t}B_k^{(i)}(x_0)\,U_{(i)\,k}^{(+)}(x-x_0)$$
$$+e^{i\varepsilon_k^{(i)}t}D_k^{(i)\dagger}(x_0)\,V_{(i)\,k}^{(-)}(x-x_0)\}\,, \tag{28}$$

$U_{(i)\,k}^{(+)}$ being a positive energy solution of the equation (21), and $V_{(i)\,k}^{(-)}$ the charge conjugate of the negative energy solution,

$$V_{(i)\,k}^{(-)}=\sigma^3\,[U_{(i)\,k}^{(-)}]^* \tag{29}$$

The operators $B_k^{(i)\dagger}(x_0)$ $(B_k^{(i)}(x_0))$ and $D_k^{(i)\dagger}(x_0)$ $(D_k^{(i)}(x_0))$ create (annihilate) conventioal quarks and anti-quarks with energy $\varepsilon_k^{(i)}$'s. $a^{(i)}{}_{\dagger}(x^0)$ $(a^{(i)}(x_0))$ creates (annihilates) a quark-soliton being localized at x_0. We assume the commutation relations

$$\begin{aligned}
&[B_k^{(i)}(x_0),\ B_{k'}^{(j)\dagger}(x_0)]_+=\delta_{ij}\,\delta(k-k')\,,\\
&[D_k^{(i)}(x_0),\ D_{k'}^{(j)\dagger}(x_0)]_+=\delta_{ij}\,\delta(k-k')\,,\\
&[a^{(i)}(x_0),a^{(j)\dagger}(x_0)]_+=\delta_{ij}\,,
\end{aligned} \tag{30}$$

and all other anti-commutators are zero. Specifically, any combination of two operators specified by x_0 and x_0' $(x_0\neq x_0')$ always anti-commutes.

The charge operator in the one soliton sector is defined by

$$Q=\int dx:\,q'^{\dagger}(x-x_0)\left(\frac{\lambda_3'}{2}+\frac{\lambda_8'}{2\sqrt{3}}\right)q':\quad, \tag{31}$$

where $\lambda_3'=U\lambda_3U^{\dagger}$ and $\lambda_8'=U\lambda_8U^{\dagger}$.
By substituting the expansion (28) into (31), we find

$$Q=\frac{2}{3}a^{(1)\dagger}(x_0)\,a^{(1)}(x_0)-\frac{1}{3}a^{(2)\dagger}(x_0)\,a^{(2)}(x_0)-\frac{1}{3}a^{(3)\dagger}(x_0)\,a^{(3)}(x_0)$$
$$+\sum_k\Big[\frac{2}{3}(B_k^{(1)\dagger}(x_0)\,B_k^{(1)}(x_0)-D_k^{(1)\dagger}(x_0)\,D_k^{(1)}(x_0))$$

We consider the non-leptonic Lagrangian

$$L_{N-L}=2\sqrt{2}\,G\sin\theta_c\cos\theta_c\,(\bar{u}\gamma_\alpha s)\,(\bar{d}\gamma^\alpha u)\,, \tag{36)*}$$

operating on the $S(x_0)$ soliton state produces another state of the same energy; hence we have many degenerate states. To distinguish them, we may lavel them as $|0;S(x_0)\rangle$ (unoccupied quark-soliton state) and $|q_{(0)i};S(x_0)\rangle$ (occupied the i-th quark-soliton state)

$$\begin{aligned} &a^{(i)}(x_0)|q_{(0)i};S(x_0)\rangle = |0;S(x_0)\rangle \\ &a^{(i)\dagger}(x_0)|0;S(x_0)\rangle = |q_{(0)i};S(x_0)\rangle \\ &a^{(i)}(x_0)|q_{(0)i};S(x_0)\rangle = 0 \\ &a^{(i)\dagger}(x_0)|q_{(0)i};S(x_0)\rangle = 0 \\ &\qquad\qquad \text{etc.}\quad . \end{aligned} \tag{33}$$

The charge quantum number of $|0;S(x_0)\rangle$ and $|q_{(0)i};S(x_0)\rangle$ are given by

$$\begin{aligned} &Q|0;S(x_0)\rangle = 0\,, \\ &Q|q_{(0)1};S(x_0)\rangle = \frac{2}{3}\,,\quad Q|q_{(0)2};S(x_0)\rangle = -\frac{1}{3} \\ &Q|q_{(0)3};S(x_0)\rangle = -\frac{1}{3}\,. \end{aligned}$$

The charge quantum number of $|q_{(0)i};S(x_0)\rangle$, therefore, is the same as the ordinary quark state.

4. SOLITON CONTRIBUTION TO NON-LEPTONIC WEAK INTERACTION

For an example of non-leptonic Lagrangian, we take the form of charge current product,

$$L = 2\sqrt{2}G\,(\bar{q}\lambda_+\gamma_\alpha q)\,(\bar{q}\lambda_-\gamma^\alpha q), \tag{34}$$

with $\lambda_+ = \frac{1}{2}(\lambda_1 + i\lambda_2)$, $\lambda_- = \frac{1}{2}(\lambda_1 - i\lambda_2)$.

The unitary transformation (14) reduces it to

$$\begin{aligned} L' &= 2\sqrt{2}G\,(\bar{q}'U\lambda_+U^\dagger\gamma_\alpha q')\,(\bar{q}'U\lambda_-U^\dagger\gamma^\alpha q') \\ &= 2\sqrt{2}G\,[\cos^2\theta_c\,(\bar{u}'\gamma_\alpha d')\,(\bar{d}'\gamma^\alpha u') \\ &\quad - e^{i\varphi}\sin\theta_c\cos\theta_c\,(\bar{u}'\gamma_\alpha s')\,(\bar{d}'\gamma^\alpha u') \\ &\quad - e^{-i\varphi}\sin\theta_c\cos\theta_c\,(\bar{u}'\gamma_\alpha d')\,(\bar{s}'\gamma^\alpha u') \\ &\quad + \sin^2\theta_c\,(\bar{u}'\gamma_\alpha s')\,(\bar{s}'\gamma^\alpha u')\,]\,. \end{aligned} \tag{35}$$

The second term (the third term) of this expression transforms a s'-quark to a d'-quark (a d'-quark to a s'-quark). We see that the initial current-current product of (34) leads to the non-leptonic Lagrangian of (34) in our SU(3)-broken symmetry.

Hereafter, we discuss the process which is caused by the second term of (35). For the sake of brevity, omitting the prime notation of q', let us represent the transformed quark by $q \equiv (u, d, s)$ in the following.

We consider the non-leptonic Lagrangian

$$L_{N-L}=2\sqrt{2}\,G\sin\theta_c\cos\theta_c\,(\bar{u}\gamma_\alpha s)\,(\bar{d}\gamma^\alpha u)\,, \tag{36}$$ *

which gives rise to transition us→ud (see Fig. 3).

If we take the weak boson and the color gluon interactions for the non-leptonic Lagrangian in a more realistic model, the bare Fermi vertex (Fig. 3) would be replaced by those graphs in which a W-boson connects two different quark lines and the vertices are dressed by gluons.[3]

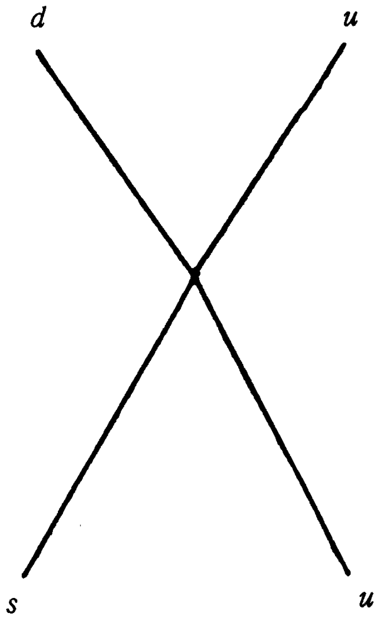

Fig. 3. Fermi vertex for the transition us→ud.

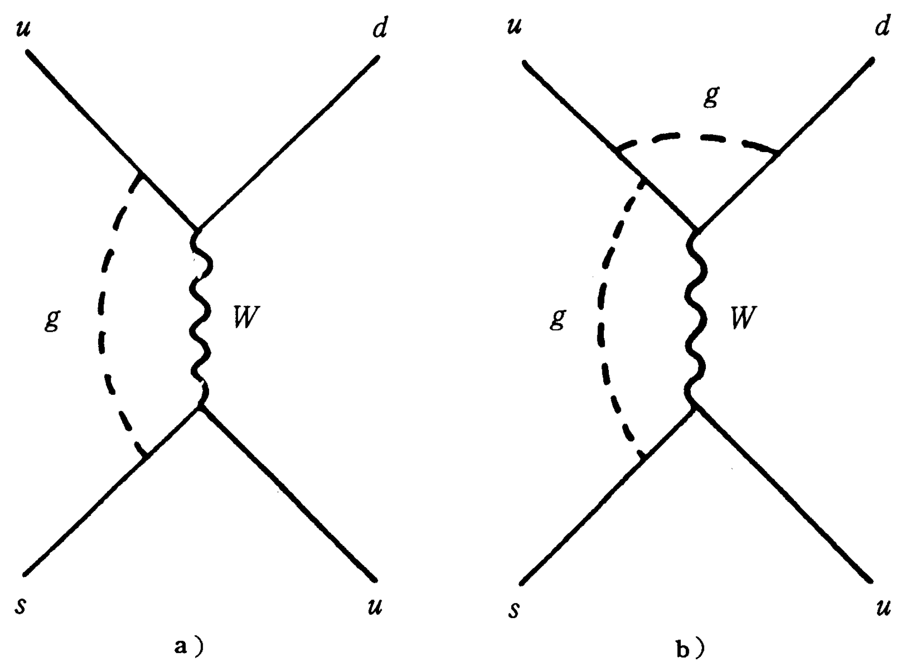

Fig. 4. g repesents a gluon and W is a weak boson.

* In a case of three spatial dimension, u, d and s are replaced with the left-handedd u_L, d_L and s_L in (36).

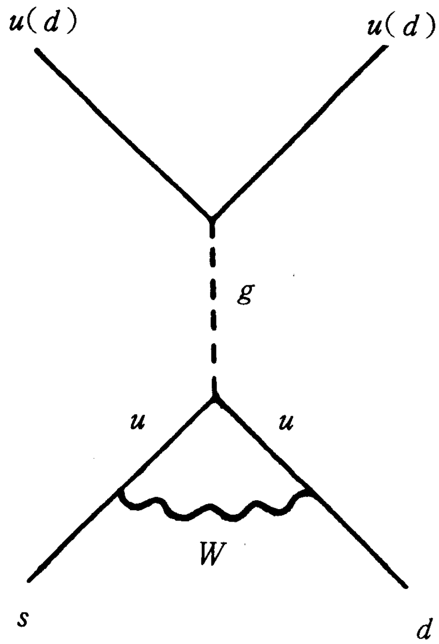

Fig. 5. Gluonic monopole graph.

Some examples of such elementary vertices is shown in Fig. 4 and 5 where the color indices of quarks and gluons are omitted.

It was known that in the standard model of three spatial dimensions, the types of graphs ot Fig. 4 contribute to the renormalization of weak coupling constant; the enhancement of the $\triangle T=\frac{1}{2}$ transition and the supression of the $\triangle T=\frac{3}{2}$ transition.[4] It also was demonstrated that, taking account that m_W(W-boson mass) $\gg m_q$(quark mass), the graph of Fig. 5 can be replaced with the product of a four-quark point

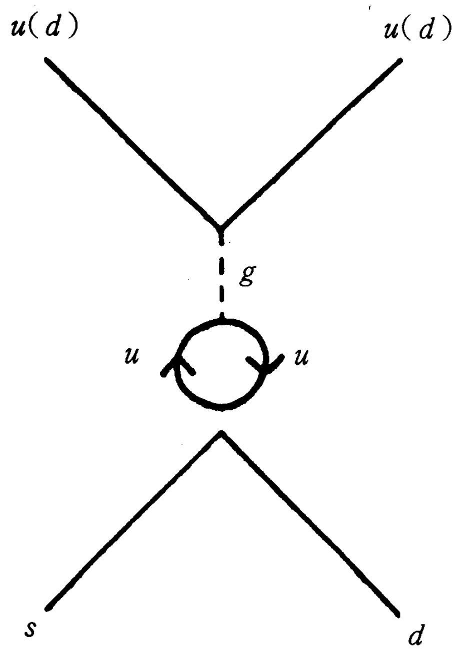

Fig. 6. Reduced gluonic monopole graph.

interaction and a gluonic monopole one shown in Fig. 6, which is expected, as a crude approximation, to enhance the right-handed quark contribution to the non-leptonic hyperon decays.[5]

Inspecting the color exchange or non-exchange character of each vertex in Fig. 4 and 6, we can see that the quark-soliton states could affect the lower vertex in Fig. 6, as we are assuming a white soliton. In order to estimate the soliton-state contributions to the non-leptonic vertex in our one-spatial model, at the first, we consider the transition of an ordinary quark to a quark-soliton in the soliton field. The vertex corresponding to this process is shown by Fig. 7.

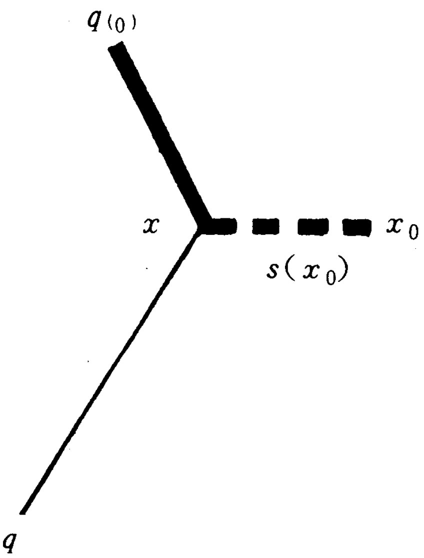

Fig. 7. Soliton and quark interaction vertex.
The bold line represents a soliton-quark line and the bold broken line shows an extension of soliton.

The transition matrix element ot an ordinary quark (q) to a quark-soliton ($q_{(0)}$) in a presence of soliton field ($S(x_0)$) is given by

$$\begin{aligned} H(k\,;x_0) &= \langle q_{(0)}\,; S(x_0) \,|\, g\mu\bar{q}\lambda_\alpha q\phi_\alpha \,|\, q_i(k)\rangle \\ &= g\mu C_i \int dx\, e^{-m_0|x-x_0|}\{\cosh\lambda(x-x_0)\}^{-g_i\mu/\lambda} \\ &\quad \times \tanh\lambda(x-x_0)\, e^{ikx}, \end{aligned} \tag{37}$$

$|q_{(0)i}; S(x_0)\rangle$ being a quark-soliton state.

Assuming that a quark-soliton wave function $C_i e^{-m_0|x-x_0|}\{\cosh\lambda(x-x_0)\}^{-g_i\mu/\lambda}$ is localized in an approximate range of $\dfrac{1}{m_0}$ around x_0, we can estimate the x-integrals in (37) and (24), then we have

$$H(k\,;x_0) \simeq ig\mu\frac{4\lambda k}{{m_0}^3} C_i e^{ikx_0} \tag{38}$$

with $C_i \simeq \sqrt{m_0}$.

Let us consider the $s(k)-d(k')$ inelastic scattering into the state $s_{(0)}-d_{(0)}$ as shown in Fig. 8.

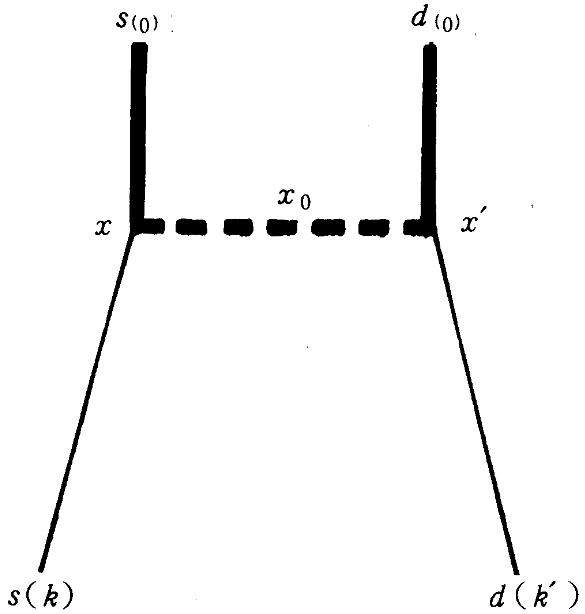

Fig. 8. s-d inelastic scattering into the $s_{(0)}$-$d_{(0)}$ state. k(k′) is a s(d)-momentum.

Making use of the sudden approximation method, the scattering amplitude is defined by

$$b\delta(k+k') = (-i)^2 2\pi \frac{m_0\, m_0}{\lambda k\, \lambda k'} \int dx_0 H_3(k\,;x_0)\, H_2(k'\,;x_0), \tag{39}$$

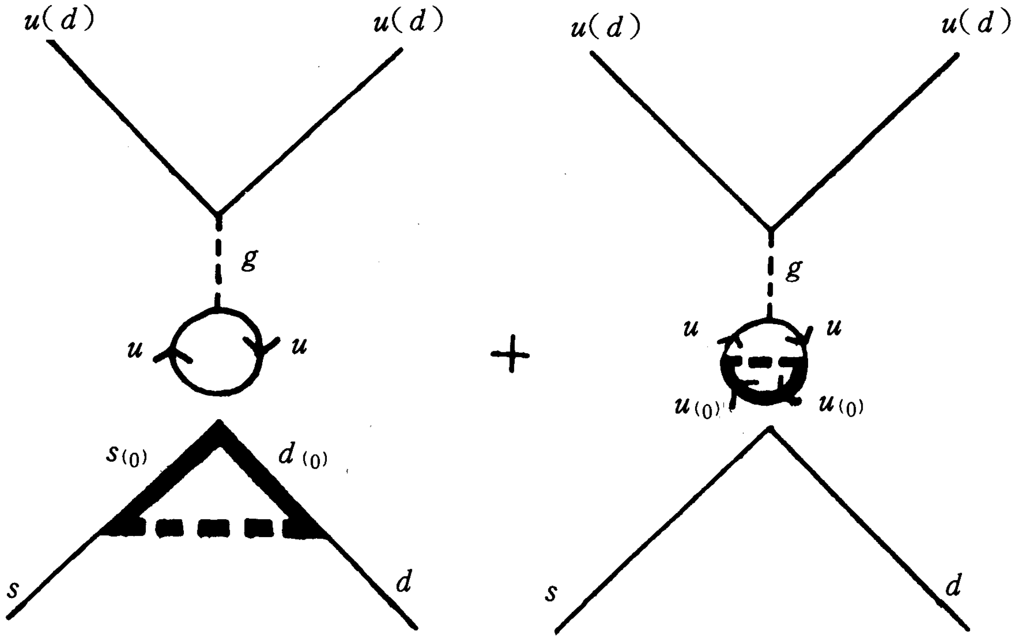

Fig. 9. The quark-soliton contribution to the gluonic monopole vertex.

where H_3 and H_2 are defined by (37) and $\frac{m_0}{\lambda k}\left(\frac{m_0}{\lambda k'}\right)$ is a $s(d)$ quark travelling time through the range $\frac{1}{\lambda}$ of soliton field. Substituting (38) into (39), we obtain

$$b \simeq \frac{32\pi g^2\mu^2 C_2 C_3}{m_0{}^4} = 32\pi \frac{g^2\mu^2}{m_0{}^3}. \tag{40}$$

The additional graph to Fig. 6 is shown in Fig. 9. In this process, the Fermi coupling constant G of four-quark vertex is replaced by bG. If we take $g\mu/m_0 \simeq 10^{-1}$, and $m_0 \simeq 30$Mev, we obtain

$$b \simeq 10^{-10} \tag{41}$$

We, therefore, conclude that the soliton state contributions to the non-leptonic hyperon-decay are negligible small.

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