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ON THE SENSITIVITY OF LOW ENERGY
WEAK INTERACTIONS TO LARGE MASS SCALES

I.I. Bigi
Inst. of Theoret. Physics E, RWTH Aachen
D5100 Aachen, Fed. Rep. Germany

and

TRIUMF
4004 Wesbrook Mall, Vancouver, B.C., Canada V6T 2A3

Abstract

It is discussed how known phenomenology from kaon, hyperon, charm and bottom decays makes it unlikely (though not impossible) that new states as suggested by composite models — namely spin zero partners of W bosons, lepto-quarks or colour octet bosons — exist with a mass less than 1 TeV.

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I. Introduction. Although the Standard $SU(3)_C \times SU(2)_L \times U(1)$ gauge model of particle physics has scored many striking successes, it is not considered to be the final theory. Numerous theoretical frameworks have been invoked to remedy some of the alleged deficiencies of the Standard Model. They hold out the promise of being more natural and economical in principle, yet at present their trademark lies in the introduction of a host of additional quanta with new spin, isospin and color assignments.

In this note we will analyze to which extent such new quanta would make their presence felt in the phenomenology of weak interactions at low energies. To present a somewhat more specific discussion we will concentrate on composite models [1] and the spectrum of new quanta they suggest — yet the results obtained are of a more general nature:

(0) Higher excitations would produce heavier partners of the W and Z^0 bosons, both with isospin one and zero. The mass of the vectorbosons typically has to exceed 1 TeV in order not to upset the observed size of the Weinberg angle. We will not comment any further on it.

(1) Scalar partners of $W^\pm$ and Z^0 should exist with isospin one — $T^\pm$, T^0 — and zero — V^0 . There is no good reason why these should be very heavy. Indeed the analogy with QCD or mass inequalities derived for vector like theories suggest

$$M_T < M_W$$

(ii) All the states listed above are singlets under $SU(3)_C$; in addition there could be colour octet quanta of spin and isospin one and zero. Furthermore there might be leptoquarks LQ, i.e.

bosons with both lepton and quark quantum numbers. Again, a priori they need not be very heavy.

Since there is no intrinsic reason why scalar bosons T, leptoquarks LQ and colour octet bosons $\tilde{W}$ are much heavier than W and Z, the exciting prospect appears that they can be produced at existing accelerators like the CERN $p\bar{p}$ collider.

The topic of this paper is to show that the known phenomenology of low-energy weak processes places very non-trivial constraints on the masses and couplings of such "new" quanta: it appears unlikely that their masses will lie below 200 GeV — actually 1 TeV and above is a much more typical scale. These constraints are derived from some of the most striking findings of low energy weak phenomenology: the helicity suppression in $\pi^+ \rightarrow \ell^+ \nu_\ell$ decays; the $\Delta I=1/2$ rule in hyperon decays; the K_L - K_S mass difference; the strong suppression of flavour changing neutral currents in K decays. In addition we will rely on an analysis of a less "classical" field, namely exclusive two-body decays of charm and bottom hadrons together with some observations from inclusive decays.

It is a sad, yet well-known fact that one is not able to calculate very reliably the various decay rates referred to above. Thus it would appear a frivolous undertaking to derive useful constraints on new dynamical elements from such processes — were it not for the fact that transition rates due to the exchange of new quanta often get enhanced relative to the conventional rates. These enhancement factors can be traced back to:

— the different chirality structure of the new effective operators; this tends to increase significantly the size of QCD radiative

corrections and can at the same time lead to much larger matrix elements;

— a different colour flow in the new operators again producing larger matrix elements.

II. Scalar Bosons (Colour Singlets): (a) $\pi \rightarrow \ell \nu$ decays. The

apparent helicity suppression in these decays puts severe constraints on the couplings of the scalar bosons $T^\pm$ to quarks and leptons. Requiring such T contributions not to exceed the radiative corrections calculated in the Standard Model one finds [2]

$$\frac{g_T^f M_T^2}{g_A^f M_A^2} < 5 \times 10^{-6} \quad (2.1)$$

Thus, if $g_T^f \sim g_A^f$, one obtains $M_T \gtrsim 30$ TeV! One can however argue that $g_T^{ff'}$, i.e. the coupling of $T^\pm$ to the fermions f and f' has to be determined by m_f or $m_{f'}/M_T$ since scalar couplings violate chiral invariance. In that case no useful limit of M_T results from (2.1).

(b) $\Delta I=3/2$ hyperon decays. T exchange leads to an effective interaction for $\Delta S=1$ transitions which will in general contain the piece:

$$L_{T,eff}(\Delta S=1) = \frac{g_T^2}{(-M_T^2)} (\bar{u}_R s_L)(\bar{d}_L u_R) + h.c. \quad (2.2)$$

This operator is multiplicatively renormalized by QCD radiative corrections; the leading log approximation yields [3]

$$L_{T,eff}^{QCD}(\Delta S=1) = \frac{g_T^2}{(-M_T^2)} c_1 (\bar{u}_R s_L)(\bar{d}_L u_R) + h.c. \quad (2.3)$$

with $c_1 = [a_S(\mu)/a_S(M_T)]^{\gamma_1}$, $\gamma_1=8/b$, $b=11-2/3$ nf and μ being a hadronic scale of order 500 MeV. One then finds

$$c_1 \sim 5 \quad (2.4)$$

1.e. a large enhancement. One should note [4] that this effective operator produces both $\Delta I=1/2^-$ and $3/2^-$ transitions, in marked contrast to the usual operator based on W_L exchange: O_+ which contains $\Delta I=1/2^-$ and $3/2^-$ pieces is mildly suppressed — $c_+ \sim 0.7$ — whereas the pure $\Delta I=1/2$ operators O_- is enhanced:

$$L_T^{\text{QCD}}(\Delta S=1, \Delta I=3/2) = \frac{c_+}{\sqrt{3}} \frac{g_T^2}{(-M_W^2)} (u_R s_L)(\bar{d}_L u_R) + \text{h.c.}$$

$$L_W^{\text{QCD}}(\Delta S=1, \Delta I=3/2) = \frac{c_+}{\sqrt{6}} \frac{g_T^2}{(-M_W^2)} [(\bar{u}_L \gamma_\mu s_L)(\bar{d}_L \gamma_\mu u_L) + (\bar{d}_L \gamma_\mu s_L)(\bar{u}_L \gamma_\mu u_L)] + \text{h.c.} \quad (2.5)$$

In addition the basic observation is to be made [4] that soft pions have larger matrix elements with pseudoscalar than with axialvector currents. For hyperon decays

$$\Lambda \rightarrow N\pi$$

one finds in the valence quark approximation

$$R = \frac{\langle N\pi | L_T^{\text{QCD}}(\Delta I=3/2) | \Lambda \rangle}{\langle N\pi | L_W^{\text{QCD}}(\Delta I=3/2) | \Lambda \rangle} = \frac{c_+}{c_+} \frac{m^2}{m_s m_q} \frac{g_T^2}{g_W^2} \frac{M_W^2}{M_T^2} \sim 140 \frac{g_T^2}{g_W^2} \frac{M_W^2}{M_T^2} \quad (2.6)$$

Since $\Delta I=3/2$ hyperon decays are reasonably well described by conventional dynamics — at least within a factor of two — we require

$$R < 1 \quad (2.7)$$

For $g_T^2 \sim g_W^2$ this implies

$$M_T > 1 \text{ TeV} \quad (2.8)$$

As before this bound is obviously vitiated if T couplings to fermions are determined by m_f/M_W .

(c) K_L-K_S mass difference. It has become very popular to compute Δm_K via the well-known quark box diagram. Although no convincing argument has been presented why the box contribution should saturate the experimental value, it is nevertheless estimated [5] to yield a fair approximation. Contributions due to "new" physics should therefore not be much larger than it.

The usual quark box contributions are suppressed by GIM factors; for two families they are just $(m_c/M_W)^2$. Thus new terms containing couplings proportional to fermion masses can become competitive.

(a) $L_{\text{eff}}(\Delta S=2)$. With two families (which is sufficient for our analysis) the two leading diagrams involving T exchange are shown in fig. 1. For

$$g_{Tcf} = \frac{m_c}{M_W} \cdot r \quad (2.9)$$

r being a number of order one, one obtains after integrating out the loop

$$L_{WT}^{\text{box}}(\Delta S=2) = B \bar{s}(1+\gamma_5) d \bar{\bar{s}}(1-\gamma_5) d$$

$$B = \frac{r^2 G_F^2 M_W^2}{2\pi^2 g_W^2} \frac{m_c^2}{M_W^2 - M_T^2} \log \frac{M_W^2}{M_T^2} \times \sin^2 \theta_c \cos^2 \theta_c \quad (2.10)$$

For definiteness we have assumed here that T bosons couple to the same quark states as W bosons; diagonalizing the quark mass matrix will then lead to the appearance of the same mixing angles.

Eq. (2.10) has to be compared to the standard expression

$$L_{WW}^{\text{box}}(\Delta S=2) = A \bar{s} \gamma_\mu (1-\gamma_5) d \bar{\bar{s}} \gamma_\mu (1-\gamma_5) d$$

$$A \approx \frac{G_F^2 M_W^2}{4\pi^2} \frac{m_c^2}{M_W^2} \sin^2 \theta_c \cos^2 \theta_c \quad (2.11)$$

Next one calculates the QCD radiative corrections. Since the chirality structure of L_{WT}^{box} differs quite clearly from that of L_{WT}^{box} , one has to deal with operators of different anomalous dimensions. The multiplicatively renormalized operators are

$$O_{\pm} = \bar{s}_L \gamma_{\mu} u_L u_L \gamma_{\mu} d_L \pm \bar{s}_L \gamma_{\mu} d_L u_L \gamma_{\mu} u_L \quad (2.12)$$

for W exchange and

$$O_1 = \bar{s}_R u_L \bar{u}_L d_R \quad (2.13)$$

for T exchange. The QCD radiative corrections can be lumped into short distance coefficients c_1 :

$$O_W^{\text{QCD}} = \frac{c_+}{2} O_+ + \frac{c_-}{2} O_- \quad (2.14)$$

$$O_T^{\text{QCD}} = c_1 O_1 \quad (2.15)$$

The c_1 depend on the anomalous dimensions γ_1

$$c_1 = \left[\frac{\alpha_s(\mu)}{\alpha_s(M_W)} \right]^{\gamma_1} , \quad \gamma_1 = \frac{8}{b} = 2\gamma_- = -4\gamma_+, \quad b = 11 - \frac{2}{3} n_f \quad (2.16)$$

The parameter μ denotes a subtraction point of the operator.

Describing gluons in the Landau gauge for convenience one finds just the three types of diagrams as shown in fig. 2 contributing in the leading log approximation. The QCD corrections can then be lumped together into two factors

$$L_{WT}^{\text{QCD}}(\Delta S) \approx B \left\{ \frac{1}{2} (n_{+,1} n_{-,1}) \bar{s}(1+\gamma_5) d \bar{s}(1-\gamma_5) d \right. \\ \left. + \frac{1}{4} (n_{+,1} n_{-,1}) \bar{s} \gamma_{\mu} (1+\gamma_5) d \bar{s} \gamma_{\mu} (1-\gamma_5) d + \text{h.c.} \right\} \quad (2.17)$$

The operator $\bar{s} \gamma_{\mu} (1+\gamma_5) d \bar{s} \gamma_{\mu} (1-\gamma_5) d$ is generated by the QCD corrections and therefore vanishes from L_{WT} when $n_{+,1} = n_{-,1}$. The short distance coefficients are given by

$$\eta_{\pm,1} = \frac{M_W^2 - M_T^2}{\log \frac{M_W^2}{M_T^2}} \int \frac{dk^2}{(k^2 + M_W^2)(k^2 + M_T^2)} \left[\frac{\alpha_s(k^2)}{\alpha_s(M_W^2)} \right]^{\gamma_{\pm}} \gamma_1 \left[\frac{\alpha_s(k^2)}{\alpha_s(\mu^2)} \right]^{2\gamma_{\mu}} \quad (2.18)$$

where $2\gamma_{\mu} = -8/b$.

For $M_W \sim M_T$ one obtains a very simple result

$$\eta_{+,1} \approx \eta_{-,1} \approx \left[\frac{\alpha_s(\mu^2)}{\alpha_s(M_W^2)} \right]^{8/b} \approx 4 - 7.6 \quad (2.19)$$

when Λ_{QCD} varies between 100 and 250 MeV.

For $M_W \ll M_T$ on the other hand (we anticipate that the present analysis rules against a relation $M_W \gg M_T$) one finds

$$\eta_{+,1} \approx \frac{b}{2+b} \frac{\log \frac{M_T^2}{\Lambda^2} [\alpha_s(M_T^2)]^{2/b} [\alpha_s(\mu^2)]^{8/b}}{\log \frac{M_T^2}{M_W^2} [\alpha_s(M_W^2)]^{10/b}} \left(1 - \frac{[\alpha_s(M_T)]}{[\alpha_s(M_W)]} \right)^{(b+2)/b} \\ \eta_{-,1} \approx \frac{b}{b-4} \frac{\log \frac{M_T^2}{\Lambda^2} [\alpha_s(\mu^2)]^{8/b}}{\log \frac{M_T^2}{M_W^2} [\alpha_s(M_W^2)]^{4/b}} \left(1 - \frac{[\alpha_s(M_T)]}{[\alpha_s(M_W)]} \right)^{(b-4)/b} \quad (2.20)$$

Choosing $M_W^2 \sim 25 M_T^2$ will serve to illustrate the size of Eq. (2.20):

$$\eta_{+,1} \sim 7.8 \\ \eta_{-,1} \sim 7.2 \quad (2.21)$$

i.e., in both cases $\eta_{\pm,1}$ represent rather sizeable numbers which are similar in magnitude.

(B) $\langle K^0 | L_{\text{eff}}(\Delta S=2) | \bar{K}^0 \rangle$. One encounters a PCAC enhancement factor $\frac{m_K^2}{m_S^2}$ in $\langle K^0 | L_{WT}^{\text{box}} | \bar{K}^0 \rangle$ relative to $\langle K^0 | L_{WW}^{\text{box}} | \bar{K}^0 \rangle$; m_S denotes the current mass of the strange quark. Putting everything together one finally finds

$$R = \frac{|\Delta m_K(TW)|}{|\Delta m_K(WW)|} \sim \frac{2r^2 m_K^2}{g_W^2 m_S^2} \frac{M_W^2}{M_K^2 - M_K^2} \log \frac{M_W^2}{M_K^2} \frac{\eta_{+1} + \eta_{-1}}{2\eta_{WW}^{\text{QCD}}} \quad (2.22)$$

$$\sim \frac{200r^2}{g_W^2} \frac{M_W^2}{M_W^2 - M_K^2} \log \frac{M_W^2}{M_K^2}$$

Requiring $R < 2$, i.e. that T exchange contributions are not much larger than those due to W exchange leads again to

$$M_T \gtrsim 1 \text{ TeV} . \quad (2.23)$$

We had assumed in our analysis leading to the lower limit (2.23) that T boson couplings to fermions are determined just by the masses of the latter. To avoid such unpleasantly high masses for T bosons one has to insist that these mass related couplings follow a very peculiar pattern, namely

$$g_{Tf_R f_L} \approx \frac{m_f}{M_W} , \quad g_{Tf_L f_R} \approx \frac{m_f}{M_W}$$

In that case there is a GIM-like factor $m_S m_d / M_W^2$ in L_{TW}^{box} instead of m_C^2 / M_W^2 and $M_T \sim M_W$ is still allowed by the K_L - K_S mass difference. Note that such couplings are those of isodoublet Higgs scalars — yet T are isovector states.

III. Colour-Non-Singlet Bosons. On very general though not very rigorously defined grounds one expects various types of bosonic

excitations that carry non-trivial colour quantum numbers. The presently popular composite models actually fall into two categories, namely whether they contain just colour octet bosons $\tilde{W}$ and $\tilde{Z}$ besides the colour singlet W and Z or whether they contain in addition lepto-quarks, i.e. bosons that form colour triplets

(a) Lepto-quarks. Lepto-quarks would in general produce transitions $K_L \rightarrow \mu^+ \mu^-$, $\mu^+ e^-$ and $K^+ \rightarrow \pi^+ \nu \bar{\nu}$, $\pi^+ e^- \mu^+$. Since the contributions to "new physics" cannot exceed the experimental numbers, it was concluded [6] that vector lepto-quarks have to be very massive:

$$M(\text{spin-1 LQ}) \gtrsim 60 \text{ TeV} \quad (3.1)$$

Scalar lepto-quarks are expected to have much reduced couplings to fermions since they violate chiral invariance. Assuming

$$g_{LQ \text{ ql}} \sim 0 \left(\frac{\sqrt{m_b m_l}}{250 \text{ GeV}} \right) \quad (3.2)$$

it was found [6]

$$M(\text{spin-0 LQ}) \gtrsim 1 \text{ TeV} \quad (3.3)$$

The authors of Ref. 7 replaced (3.2) by

$$g_{LQ \text{ ql}} \sim 0 \left(\frac{m_q}{250 \text{ GeV}} \right) \quad (3.2')$$

to allow for much lighter lepto-quarks:

$$M(\text{spin-0 LQ}) \gtrsim 20 \text{ GeV} \quad (3.3')$$

While known phenomenology does not rule out such a scenario, it does not appear to be a very probable one: Although light scalar lepto-quarks with a mass $\sim 20 \text{ GeV}$ might be allowed in models with "nearly

compositeness" [7], the spin-1 lepto-quarks would then appear excessively heavy (see Eq. (3.1)).

(b) Colour-Octet Bosons. (a) Inclusive decays. $\tilde{W}$ exchange produces only non-leptonic transitions $Q + \bar{q}q_2 q_1$ described by

$$L_{\text{eff}}(\tilde{W}) = \frac{g_W^2}{(-M_W^2)} \left(\bar{q}_L \gamma_\mu \frac{\lambda_1}{2} q_L \right) \left(\bar{q}_1 L \gamma_\mu \frac{\lambda_1}{2} q_2 L \right) = \frac{g_W^2}{-M_W^2} O_8 \quad (3.4)$$

O_8 can be expressed in terms of the multiplicatively renormalized operators O_i :

$$O_8 = \frac{1}{3} (O_+ - 2O_-) \quad (3.5)$$

Thus the colour factor for the inclusive decay $Q + \bar{q}q_2 q_1$ is for $\tilde{W}[W]$ exchange:

$$\frac{1}{\tilde{W}} [L_W] = \frac{2}{9} C_+^2 + \frac{4}{9} C_-^2 [2C_+^2 + C_-^2] \quad (3.6)$$

Therefore $\tilde{W}$ contributions to inclusive decays suffer a sizeable suppression relative to W terms

$$\frac{1}{\tilde{W}} \approx \frac{2}{9} [0.3] \quad \text{for } c_+ = c_- = 1 [c_+ = 0.8, c_- = 1.5] \quad (3.7)$$

Hence the branching ratio for semi-leptonic B decays gets reduced from its conventional value of 14% ($\pm 1\%$) only to 12% even if $\frac{g_W^2}{M_W^2} = g_W^2/M_W^2$.

(b) $B + \psi K$ decays. $B + \psi K$ decays are suppressed by a colour factor $f_W = [1/3(2c_+ - c_-)]^2$ in W exchange when soft gluon exchanges are ignored:

$$f_W = \frac{1}{9} [0.002] \quad \text{for } c_+ = c_- = 1 [c_+ = 0.8, c_- = 1.5] \quad (3.8)$$

The present experimental upper limit is

$$BR(B + \psi K) < 1.6\% \quad (90\% \text{ C.L.}) \quad (3.9)$$

which requires a colour suppression factor $(f_W)_{\text{exp}} \lesssim 1/9$ [8].

The corresponding factor for $\tilde{W}$ is

$$f_{\tilde{W}} = \left[\frac{2}{9} (c_+ + c_-) \right]^2 = \begin{cases} \frac{16}{81} & c_+ = c_- = 1 \\ 0.26 & c_+ = 0.8, c_- = 1.5 \end{cases} \quad \text{for} \quad (3.10)$$

Thus $\tilde{W}$ contributions to $B + \psi K$ decays are enhanced over the standard W terms due to colour effects by $f_W/f_{\tilde{W}}$ and very crudely one can conclude from (3.9):

$$\frac{g_W^2}{M_W^2} \sim \frac{1}{2} \frac{g_W^2}{M_W^2} \quad (3.11)$$

(γ) Exclusive D decays. Bauer and Stech [9] have shown that exclusive two-body decays of charm mesons can be described surprisingly well in a simple quark decay picture when final state interaction are included. Parametrizing the effective interaction by

$$L_{\text{eff}}(AC=1) = \frac{G}{\sqrt{2}} \cos^2 \theta (a_1 \bar{s}_L \gamma_\mu c_L u_L \gamma_\mu d_L + a_2 \bar{u}_L \gamma_\mu c_L s_L \gamma_\mu d_L) \quad (3.12)$$

they found that the whole observed pattern of two-body charm decays is reproduced quite well when one chooses

$$a_1 \sim 1.2, a_2 \sim -0.7 \quad (3.13)$$

In the Standard Model one finds

$$a_1 = \frac{1}{2} (c_+ + c_-) + \frac{\xi}{2} (c_+ - c_-) \quad a_2 = \frac{1}{2} (c_+ - c_-) + \frac{\xi}{2} (c_+ + c_-) \quad (3.14)$$

Here ξ is the colour suppression factor; ignoring soft gluon effects (which might be an oversimplification for charm decays) one would estimate $\xi=1/3$. For $m_c=1.5$ GeV one gets

$$a_1^{\text{th}} \approx 1.3 - 0.6 \xi \quad (3.15a)$$

$$a_2^{\text{th}} \approx -0.6 + 1.3 \xi \quad (3.15b)$$

To reconcile (3.15), in particular (3.15b) with (3.13) one has to use $\xi=0$; yet considering the uncertainties involved, a value $\xi \sim 1/3$ is not ruled out.

Including $\tilde{W}$ exchange one has to replace (3.14) by

$$a_1^{\text{th}} = \frac{1}{2} (c_+ + c_-) + \frac{\xi}{2} (c_+ - c_-) + \frac{g_{\tilde{W}}^2 M_{\tilde{W}}^2}{g_W^2 M_W^2} \left\{ \frac{c_+ - 2c_-}{3} + \xi \frac{c_+ + c_-}{3} \right\} \quad (3.16a)$$

$$a_2^{\text{th}} = \frac{1}{2} (c_+ - c_-) + \frac{\xi}{2} (c_+ + c_-) + \frac{g_{\tilde{W}}^2 M_{\tilde{W}}^2}{g_W^2 M_W^2} \left\{ \frac{c_+ + 2c_-}{3} + \xi \frac{c_+ - 2c_-}{3} \right\} \quad (3.16b)$$

Numerically

$$a_1^{\text{th}} \approx 1.3 - 0.6 \xi + \frac{g_{\tilde{W}}^2 M_{\tilde{W}}^2}{g_W^2 M_W^2} (-1 + 1.5\xi) \quad (3.17a)$$

$$a_2^{\text{th}} \approx -0.6 + 1.3 \xi + \frac{g_{\tilde{W}}^2 M_{\tilde{W}}^2}{g_W^2 M_W^2} (1.5 - \xi) \quad (3.17b)$$

Comparing (3.17) with (3.13) shows that the $\tilde{W}$ contributions have a strong tendency to destroy the agreement between the fitted and the calculated values for a_1 and in particular a_2 . Therefore we conclude

$$\frac{g_{\tilde{W}}^2}{M_{\tilde{W}}^2} < \frac{1}{9} \frac{g_W^2}{M_W^2} \quad (3.18)$$

(σ) Resume on $\tilde{W}$. Since the K_L - K_S mass difference does not give a more stringent bound, we use (3.18) to conclude:

$$M_{\tilde{W}} > \begin{cases} 240 \text{ GeV} & \text{if } g_{\tilde{W}}^2 \approx g_W^2 \\ & \text{if } g_{\tilde{W}}^2 \approx g_W^2 \end{cases} \quad (3.19a)$$

$$1 \text{ TeV} \quad \frac{g_{\tilde{W}}^2}{g_W^2} \approx \frac{\alpha_s}{\alpha} \quad (3.19b)$$

Since $\tilde{W}$ bosons carry colour, we find it quite likely that they command strong couplings to fermions; thus the bound (3.19b) appears more reasonable though less conservative.

For completeness it should be added that spin-zero octet bosons $\tilde{T}$ could be much lighter than W bosons. For reasons of chiral invariance $\tilde{T}$ couplings to fermions must be suppressed by $m_q/M_{\tilde{T}}$. Since the QCD corrections are rather small for $\tilde{T}$ exchange operators and also the matrix elements are of a reduced magnitude, no enhancement factors arises as it was the case for T exchange.

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FIGURE CAPTIONS

1. Box diagrams for $L_{WT}(\Delta S=2)$
2. Gluon diagrams contributing in the leading log approximation.

Fig. 1

Fig. 2