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PENGUIN LOOPS WITH CONFINED QUARK PROPAGATORS -
THE $\Delta I = \frac{1}{2}$ RULE AS A LONG DISTANCE EFFECT?

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# PENGUIN LOOPS WITH CONFINED QUARK PROPAGATORS - THE $\Delta I = \frac{1}{2}$ RULE AS A LONG DISTANCE EFFECT?

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Abstract: We calculate the $\Delta S = 1$ penguin diagram by representing the internal quark lines in the loop by bag model wave functions. Because of the involved GIM-mechanism we keep only the lowest internal quark modes in the loop, that is with quark momenta of order $m_c$ and lower. Our result depends cruically on the values of the strong coupling constant and on the quark energy of the bag model wavefunctions. With reasonable values of parameters, · · find contributions corresponding to effective penguin coefficients ~ 2 to 5 times the standard perturbative ones. Thus the theoretical value for the ratio between $\Delta I = 1/2$ and $\Delta I = 3/2$ amplitudes seems to be improved.

1 Introduction. It has been argued by Shifman, Vainshtein and Zakharov (SVZ) [1] that the penguin diagram can explain the $\Delta I = 1/2$ rule in non-leptonic strangeness changing ($\Delta S = 1$) decays. It is therefore necessary to determine the quantitative importance of this diagram. In a previous paper [2] we have discussed the standard treatment [1] of the penguin diagram. The effective Hamiltonian for non-leptonic $\Delta S = 1$ decays is the sum of four-quark operators multiplied by coefficients containing the effects of hard gluons [1,3]. The coefficients corresponding to non-penguin operators (- see Fig. 1) contain leading logarithmic effects ~ $\ln(M^2/\mu^2)$ where $\mu$ is the renormalization point and M the W-mass. The matrix elements of the effective Hamiltonian is of course independent of $\mu$, which is in principle arbitrary. However, using approximative methods for calculating matrix elements of the involved ($\mu$-dependent) four quark operators, $\mu$ is in the standard approach [1,3] chosen to be some typical hadronic mass $\lesssim$ 1 GeV. Thus the standard approach [1,3] is ambiguous. Numerically, however, the coefficients of non-penguin operators are not very sensitive for the specific choice of $\mu$, as long as $\mu$ is of order 1 GeV. Turning to the penguin operators [1] this is not the case: Due to the GIM-mechanism [4] the coefficients of penguin operators (- see Fig. 2) in the perturbative approach [1] involve the difference of two leading logarithm's originating from the u-quark and the c-quark loop, respectively

$$\ln(M^2/\mu^2) - \ln(M^2/m_c^2) = \ln(m_c^2/\mu^2) \qquad (1)$$

For a quark mass $m_c \cong 1.2$ to 1.5 GeV, the expression (1) is extremely sensitive to the choice of $\mu$, and for $\mu \sim m_c$ this "leading logarithmic" expression for the penguin diagram is not valid [2,5]. Moreover, even within the SVZ approach [1] where $\mu^2 \ll m_c^2$ is assumed, the coefficients $C_5$ and $C_6$ corresponding to the penguin operators $Q_5$ and $Q_6$ are *one or two orders of magnitude smaller* than the coefficients of non-penguin operators [3]. It was, however, argued [1] that the matrix elements of the penguin operators - which correspond to pure $\Delta I = 1/2$ transition -are enhanced due to their chiral structure. But still it seems to be necessary to increase the penguin coefficients "by hand" by a factor of order 5 to explain the $\Delta I = 1/2$ rule [6,7,8]. An increase of this order would naively correspond to a renormalization mass $\mu$ of order 10 MeV, which is completely unrealistic. Already at $\mu \cong 400$ MeV, say, some of the calculated effect could be part of the wave function.

Therefore a perturbative QCD approach seems to break down, i.e. we are in the region where the quarks feel the confinement. Thus for instance the MIT-bag model [9,10] could give a better description of the penguin loop diagram than the standard one [1].

Concerning the matrix elements of penguin operators for $K \to \pi\pi$, it has recently been pointed out [11] that these does not satisfy the general chiral and SU(3) properties. However, it has very recently been shown [12] that these problems can be solved by taking into account the momentum dependence of the matrix elements [12] and the socalled anomalous commutator terms [12, 13, 14]. The analysis also shows that the importance of penguin operators are numerically reduced [12, 14, 15].

The sensitivity of the specific choice of $\mu$ described above, and the failure of the SVZ approach [1] to give a satisfactory numerical explanation of the $\Delta I = \frac{1}{2}$ rule [6,7,8], motivates us to try another approach for the penguin diagram. In this paper we will describe the penguin diagram as a non-perturbative effect in the following sense: The quarks in the penguin loop (see Fig. 2b) will be described in terms of bag model wave functions. The penguin diagram will still be a u-quark loop minus a c-quark loop [4], and the amplitude can be written

$$T_P = T_u - T_c \qquad (2)$$

where $T_u(T_c)$ is the individual u(c) -quark loop amplitude corresponding to Fig. 2. For quark loop momenta $p_q$ bigger than some mass $m_*$ of order $m_c$ we expect that perturbative QCD can be used, and we obtain contributions $\simeq \ln(M^2/m_*^2)$ in the "leading logarithmic approximation" (LLA), both for the u-quark and c-quark loops. (Moreover, we don't think it makes sense to use the bag model for higher modes with quark energy-momenta which are clearly in the perturbative region). Thus, for quark momenta $p_q \gtrsim m_*$ we expect that the contributions from the u- and c-quark loops will roughly cancel each other, and that the total contribution is determined by quark momenta $p_q < m_*$:

$$T_P \simeq T_P(p_q < m_*) = T_u(p_u < m_*) - T_c(p_c < m_*) \qquad (3)$$

which corresponds to (1) in LLA. However, we will calculate the loop contributions $T_u(p_u < m_*)$ and $T_c(p_c < m_*)$ inside the bag where the quarks feel the confinement [≠1]. Note that if $m_* \lesssim m_c$, the whole c-quark loop contributions belong to the perturbative region, and $T_p \cong T_u$ ($p_u < m_*$). It should be emphasized that in our approach there are no explicite penguin operators as in ref. 1. Our effective Hamiltonian is that of ref. 3, with $\mu = m_*$.

We expect that our bagmodel approach to the penguin loop diagram described above is too idealized and simplified to tell the whole truth about the penguin effect. But in this paper we will consentrate on the main results of this approach. We intend to give a more systematic presentation of the details of this work combined with a discussion of various aspects of the penguin diagram in a later paper.

2. Penguin loop calculation. After a Fierz-transformation (- assuming W to be heavy -) the penguin diagram will look like Fig. 2b. The amplitude for this diagram can be written:

$$T = g_s^2 G_W^{eff} \iint d^4x\, d^4y\; J_\mu^a(x; s\to d)\, \Pi^{\sigma\mu}(x,y)\, A_\sigma^a(y;q) \quad , \tag{4}$$

where $g_s$ is the quark-gluon coupling. $G_W^{eff}$ and $J_\mu^a(x;s\to d)$ are the effective weak coupling and the transition current originating from the well-known effective Hamiltonian [3] acting at the upper vertex in Fig. 2b (space-time point x). $A_\sigma^a(y,q)$ is the standard "colour-electromagnetic" potential [10,16] due to the quark-gluon interaction (for quark flavour q = u,d,s) at the lower vertex in Fig. 2b. $\Pi^{\sigma\mu}(x,y)$ is the quark loop tensor calculated within the MIT bag model. The quantities involved in eq. (4) will be explicitely given below:

$$G_W^{eff} = \sqrt{2}\, G_F \sin\theta_c \cos\theta_c \; (C_+ + C_-) \quad , \tag{5}$$

where $G_F$ is the Fermi coupling constant, $\theta_c$ is the Cabibbo angle, and the coefficients $C_\pm$ contains the effect of hard gluons [3]. Numerically, $C_- \cong 2.5$ and $C_+ \cong 0.7$, while the free field values are $C_\pm = 1$. The (coloured) transition current in (4) is

$$J_\mu^a(x; s\to d) = \overline{\psi_d(x)}\, \gamma_\mu \tfrac{1}{2}\lambda^a L\, \psi_s(x) \quad , \tag{6}$$

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≠1 To calculate the loop diagrams for weak interactions inside the bag has recently also been suggested by Donoghue [8], whom I would like to thank for discussion at an early stage of this work.

where $L = (1-\gamma_5)/2$, $\lambda^a$ are the SU(3) Gell-Mann matrices, and $\psi_s(\psi_d)$ is the s(d) quark field. The potential $A^a_\sigma$ in (4) has the form [11]

$$A^a_\sigma(y;q) = \frac{1}{2}\lambda^a\left(A_0(\rho;q)\,,\ \hat{\rho}\times\sigma A(\rho;q)\right)\,, \quad (7)$$

where $\rho = |\chi|$. The quark loop tensor is

$$\Pi^{\sigma\mu}(x,y) = Tr\left\{S(x,y)\gamma^\sigma S(y,x)\gamma^\mu L\right\}\,, \quad (8)$$

where S(x,y) is the quark propagator in terms of bag model wave functions

$$S(x,y) = -i\sum_N \Big\{\theta(x_0-y_0)\,\psi_N(x)\,\overline{\psi_N(y)}$$

$$-\theta(y_0-x_0)\,\psi^C_N(x)\,\overline{\psi^C_N(y)}\Big\}\,. \quad (9)$$

$\psi_N$ and $\psi^C_N$ represent quark and antiquark (charge-conjugated) modes respectively (The frequency factor $\exp\{-iE_N x_0\}$ is included in $\psi_N(x)$). N symbolizes some specific mode: $1S_{1/2}$, $1P_{1/2}$, $1P_{3/2}$, $2S_{1/2}$, ... in the MIT-bag model.

The amplitude T in eq. (4) is the individual u-quark ($T_u$) or c-quark ($T_c$) loop amplitude, and the penguin contribution ($T_P$) is given by eq. (2). Both propagators S in eq. (8) are a sum over quark modes N (see eq. 9). Thus (8) is a double sum over quark modes N and N'. From the time integrations $\int\int dx_0 dy_0$ in eq. (4) we obtain a δ-function expressing conservation of energy, and an energy denominator $(E_N + E_{N'})^{-1}$. (Note the relative + sign in the energy denominator due to our particle-antiparticle loop). N and N' represent internal quark modes running at the right and left part of quark loop of Fig. 2b. Removing from T 2πi times the δ-function expressing energy conservation, we obtain the following T-matrix for $K \to \pi$ transition of the following form:

$$\langle\pi|\hat{T}|K\rangle = \sum_{N,N'} \alpha_s \frac{U(N,N')}{\omega_N + \omega_{N'}}\,, \quad (10)$$

where the U(N,N')'s contain the integrals over bag wave functions, and $\langle\lambda^a\lambda^a\rangle$ is the matrix element of $\lambda^a\lambda^a$ in colour space, which is -8/3 for baryon to baryon and +16/3 for meson to meson matrix elements respectively. Moreover, $\omega_N = RE_N$, R being the bag radius.

For given $N$ and $N'$ the $U(N,N')$'s in (10) can be written as a product $U = U_x \hat{U}_y$ where $U_x$ and $\hat{U}_y$ are bag integrals over $\int d^3x$ and $\int d^3y$ respectively. The $U_x$'s are integrals over products of four (upper and/or lower component) bag wave functions. This type of integrals always enter in standard calculations of local four quark operators in weak decays [2,7,13,17], and their values are typically of order $10^{-3} \times$ GeV$^3$.

The $\hat{U}_y$'s are integrals over a product of two bag model wave functions and the vector potential in eq. (7). Such integrals are involved in colour magnetic splitting calculations [10,16,18,19] in hadron spectroscopy. For the lowest lying bag modes they are typically of order $10^{-1}$. From (10) we observe that our result depends cruically on $\alpha_s$ and the quark energy modes $E_N = \omega_N/R$. Contributions from higher quark modes in the loop will be damped by the factor $(\omega_N + \omega_{N'})^{-1}$.

The matrix element of the penguin part of the effective Hamiltonian $H_W^P = \hat{G}(C_5Q_5 + C_6Q_6)$ for a weak transition $A \to A'$ within the SVZ approach [1] can be written

$$\langle A'|H_W^P|A\rangle = C_{eff}^P \hat{G} \langle A'|Q_6|A\rangle \qquad (11)$$

where $\hat{G} \equiv \sqrt{2}\, G_F \sin\theta_c \cos\theta_c$ and $Q_6 \equiv \bar{\psi}_d \gamma_\mu L \psi_s \bar{\psi}_q \gamma^\mu R \psi_q$. ($R \equiv (1 + \gamma_5)/2$ is the right-handed projector in Dirac-space). The effective penguin coefficient $C_{eff}^P$ is $-8/3\ C_5 + C_6$ for baryon to baryon transitions and $+16/3\ C_5 + C_6$ for meson to meson transitions. For $K \to \pi$

$$-\sqrt{2}\ \langle \pi^0|Q_6|K^0\rangle = \langle \pi^\pm|Q_6|K^\pm\rangle = a + b \quad , \qquad (12)$$

$$\frac{1}{\sqrt{2}} \langle \pi^0|Q_4|K^0\rangle = \langle \pi^\pm|Q_4|K^\pm\rangle = -2(a-3b) \quad , \qquad (13)$$

where $a$ and $b$ are defined in ref. 17. $H_w(\Delta I=3/2) = \hat{G}C_4Q_4$ is the part of the effective Hamitonian responsible for $\Delta I = 3/2$ transitions. In the "bag penguin loop" approach chosen in this paper there is no penguin effective Hamiltonian $H_W^P \to \hat{G}C_{eff}^P Q_6$. However, it is convenient to define an "effective bag penguin loop coefficient" for the weak transition $A \to A'$:

$$C_{eff}^P(\text{Bag}) \equiv \frac{\langle A'|\hat{T}_P|A\rangle}{\hat{G}\langle A'|Q_6|A\rangle} \qquad (14)$$

Note that $C^P_{eff}$ is obtained if $f_P$ is replaced by $K^P_V \cdot C^P_{eff}$(Bag) is of course process dependent, while this is normally not supposed to be the case for the standard effective penguin coefficient $C^P_{eff}$ [1].

3 Variation of the strong coupling. Calculating one-gluon exchange diagrams for mass splittings within the MIT-bag model, the fine structure constant $\alpha_s = g_s^2/4\pi$ is taken to be ≈ 2.2 to fit the observed mass differences among ordinary hadrons built of the lowest lying $1S_{1/2}$ quark modes [1C]. For mass splittings among negative parity baryons containing an excited P-wave ($1P_{1/2}$ or $1P_{3/2}$) quark, the corresponding effective $\alpha_s$ seems to be smaller [18]. In fact, the best fit gives $\alpha_s \simeq 0.7$[≠2]. It should also be noted that to fit the $\Sigma - \Lambda$ mass difference, the effective $\alpha_s$ has to be reduced from the standard value (≈2.2, say) by roughly 20% when one gluon is emitted from an $1S_{1/2}$ s-quark mode [20]. *Thus one cannot use the same $\alpha_s$ for all bag model calculations.*

From perturbative QCD and the renormalization group equations the variation of $\alpha_s$ with respect to the momentum scale Q can be predicted. But a priori we cannot expect that this effective $\alpha_s$ can be used in bag model calculations for the lowest light quark modes. But we expect that the renormalization group formula for decrease in $\alpha_s$ with *quark energy is more reliable* for higher modes. Moreover, when the gluon couples to a charmed quark mode, we expect the coupling constant to be much smaller than for the corresponding u-quark modes. For instance, in charmonium decays $\alpha_s \sim 0.2$ and in the SVZ approach [1] $\alpha_s(m_c^2) \simeq 0.4$. With the above comments in mind we have chosen some reasonable values of $\alpha_s$ for the lowest modes (N, N') in the penguin loop. These are given in table 1 together with the corresponding results.

It should be noted that when a gluon is exchanged between two quark-gluon vertices 1 and 2 with coupling g(1) and g(2), the effective $\alpha_s$ is taken to be

$$\alpha_s^{eff} = \frac{g_s(1)g_s(2)}{4\pi} = \sqrt{\alpha_s(1)\,\alpha_s(2)} \qquad (15)$$

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≠2 One should note, however, that the authors of ref. 18 includes effects of the pion field outside the bag, and use $\alpha_s = 1.55$ instead of 2.2 for ordinary non-excited hadrons.

4. Results. The contributions to $C^P_{eff}(Bag)$ from the lowest bag model modes are given in table 1. The total result for $C^P_{eff}(Bag)$ depends on the choice of the momentum scale $m_*$ dividing the perturbative and the bag model region. Especially when the variation of the fine structure constant $\alpha_s^{eff}$ described in section 3 is taken into account, the lowest modes give the most sizable contribution. It should be noted that we only get non-zero contributions to the amplitude T when N(N') is an S-mode and N'(N) a P-mode. Thus the quark-antiquark pair in the penguin loop has positive parity and the bag model matrix elements for meson to meson (- and similarly for baryon to baryon) weak transitions turns out to be parity conserving as usual [2,17,22].

Our result is given in table 1. $\alpha_s(N,N')$ is the strong fine structure constant corresponding to the upper quark gluon vertex (at space-time point y) in Fig. 2. $N,N' = 1S_{1/2}, 1P_{1/2}$ .... denotes the various modes for the quarks in the penguin loop. At the lower quark gluon vertex, the gluon couples to one of the valence quarks and $\alpha_s$ is always $\alpha_s(1S_{1/2},1S_{1/2})$. The effective $\alpha_s$ for the penguin diagram for given N,N' is then given by eq. (15). The contributions to $C^P_{eff}(Bag)$ due to the lowest modes in the penguin loop are given in the table. Now, assume that $m_*$ is chosen such that the whole c-quark loop contribution belongs to the perturbative region; i.e. the quark energy-momenta $p_c$ satisfy $p_c \gtrsim m_*$ and $T_p \cong T_u(p_u < m_*)$ (- see eq. 13)). Examining numerically the u-quark energy-momenta for the MIT-bag, we find that we have to sum (N,N') = (1S,1P), (1S,2P) and (2S,1P) modes for the u-quark loop. Thus we obtain

$$C^P_{eff}(Bag) \simeq -0.9 \qquad (16)$$

which is roughly three times the standard perturbative result [1]. If $m_*$ is increased such that the lowest (1S,1P) modes for the c-quark loop are included in the non-perturbative region, we also have to include (2S,2P) modes for the u-quark loop. It is seen from table 1 that the sum of these contributions roughly cancel due to the GIM-mechanism (- see eq.(3)). Thus our result does not depend critically on $m_*$.

The value in (16) is possibly still not big enough to explain the $\Delta I = 1/2$ rule. However, if the simple MIT-version [9,10] of the bag model is modified, the quark energies tend to be smaller. Both in chiral bag models [23] and in the socalled LAPP-version [24] the quark energies are typically smaller by a factor ~1.5, and we would expect to obtain a

$C_P^{eff}$(Bag) which is 4 to 5 times the perturbative one [1]. According to refs.7,8 this could be enough to explain the $\Delta I = 1/2$ rule. However, there are still some uncertain points to consider. For instance it has been shown [14] that matrix-elements of non-penguin operators of the effective Hamiltonian $H_w$[1,3] are overestimated in quark models. In particular the matrix element for $\Delta I = 3/2$ in (13) is overestimated in bag models. Moreover, using the soft-pion theorem to relate $\langle\pi\pi|H_w|K\rangle$ to $\langle\pi|H_w|K\rangle$, the penguin operator $Q_6$ in eqs. (11)-(14) should be replaced by $Q_6^{eff} = Q_6 + Q_6^{(c)}$, where $Q_6^{(c)}$ is the socalled anomalous commutator term [12,13,14]. (This is not the case for non-penguin operators of left-left type). This will effectively reduce (12) and enhance $C^P_{eff}$(Bag) in (14). (However, the absolute value of $\langle\pi|\hat{T}_p|K\rangle$ in (10) and (14) is still the same).

We have used the standard value $\alpha_s = 2.2$ for $1S_{1/2}$ light quark modes, with reasonable smaller values for higher modes (- see table 1). Other bag models [25] use smaller $\alpha_s$'s (- and smaller R's). This will reduce the value of $C^P_{eff}$(Bag) and make the explanation of the $\Delta I = 1/2$ rule in terms of our "bag-penguin" effect more doubtful.

5. Discussion. We have calculated penguin loop contributions which are non-perturbative - in the sense that the quark propagators takes into account confinement effects in terms of the bag model quark wave functions. We have defined a momentum scale $m_*$ of order $m_c$ which divides the perturbative region and the non-perturbative region described by bag model quark wave functions. Because of the involved GIM-mechanism [4], the perturbative contributions $\sim \ln(M^2/m_*^2)$ cancels in LLA. This strict division is of course somewhat idealized, but because the lowest modes gives the dominating contributions, we find our method reasonable. It should be noted that according to our philosophy, the top quark penguin loop contribution belongs to the perturbative region. But the sizeable logarithmic contribution $\sim \ln(m_t^2/m_*^2)$ is suppressed by additional generalized Cabibbo factors.

Qualitatively, we agree with SVZ [1] that the penguin diagram could hopefully explain the $\Delta I = \frac{1}{2}$ rule for $K \to \pi\pi$ decays. However, our procedure is different because we don't include the penguin four quark operators ($Q_5$ and $Q_6$) in the effectice Hamiltonian. While a standard bag model calculation of $K \to \pi\pi$ [7,13] based on the SVZ effective Hamiltonian [1] fails to give big enough enhancement of the penguin loop contribution, it seems that our treatment based on bag model loops can account for bigger penguin contributions to $\Delta I = 1/2$ amplitudes. Concerning chiral properties, our result does not behave as the inverse of current quark masses as in Ref. 1. And momentum dependence to ensure the correct SU(3) and chiral properties of the amplitudes [11,15] can be implemented as in refs. 12,13. However, one should be careful at this point because the standard MIT-bag model does not satisfy chiral invariance. Further comments and details concerning this point will be postponed to a later paper.

Wise and Witten [24] have shown that in the standard SVZ approach to the penguin diagram, where $m_c$ is regarded to be heavy, higher order gluon effects do not spoil the one-gluon exchange SVZ result [1]. We have not explicitely considered how multigluon effects will influence our "bag-penguin" result. However, as in the case of chromomagnetic mass-splitting interactions [10,11,13], we may expect that our "bag-penguin" effect can be fairly well described as a one gluon exchange effect.

As explained in section 4, there are some numerical uncertainties in the calculations of $C_p^{eff}$(Bag) as well as the $\Delta I = 3/2$ amplitude. Moreover, our approach is simplified and the final explanation could be an interplay of the mechanism proposed here and (- or overlap with -) those proposed in previous papers [1,27,28,29,30,31]. However, despite the numerical uncertainties we think that the enhancement of the long distance penguin amplitude obtained here is encouraging and hopefully indicates that we are on the right track of the explanation of the $\Delta I = 1/2$ rule.

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Table 1. Contributions to the "effective bag penguin loop" coefficient $C_p^{eff}$(Bag) for the lowest quark modes (The minus sign due to GIM is included)

| N | N' | u-quark $\alpha_s(N,N')$ | u-quark $C^P_{eff}$(Bag) | c-quark $\alpha_s(N,N')$ | c-quark $C^P_{eff}$(Bag) |
|---|---|---|---|---|---|
| $1S_{1/2}$ | $1S_{1/2}$ | 2.2 | 0 | 0.30 | 0 |
| $1S_{1/2}$ | $1P_{1/2}$ | 1.2 | -0.26 | 0.27 | +0.13 |
| $1S_{1/2}$ | $1P_{3/2}$ | 1.2 | -0.36 | 0.27 | +0.16 |
| $2S_{1/2}$ | $1P_{1/2}$ | 1.0 | -0.14 | 0.25 | $10^{-3}$ |
| $2S_{1/2}$ | $1P_{3/2}$ | 1.0 | -0.05 | 0.25 | $3\times10^{-3}$ |
| $1S_{1/2}$ | $2P_{1/2}$ | 1.0 | $-3\times10^{-3}$ | 0.25 | 0.02 |
| $1S_{1/2}$ | $2P_{3/2}$ | 1.0 | $-4\times10^{-3}$ | 0.25 | 0.01 |
| $2S_{1/2}$ | $2P_{1/2}$ | 0.9 | -0.12 | 0.22 | 0.06 |
| $2S_{1/2}$ | $2P_{3/2}$ | 0.9 | -0.12 | 0.22 | 0.06 |

## Figure Captions

Fig. 1. Feynman diagrams corresponding to the effective Hamiltonian of Ref. 3.

a) The bare diagram.

b) QCD correction containing ~ $\ln(M^2/\mu^2)$ effects.

Fig. 2. The penguin loop diagram.

a) -for finite M.

b) - Fierz-transformed version for $M \to \infty$.

Fig.1

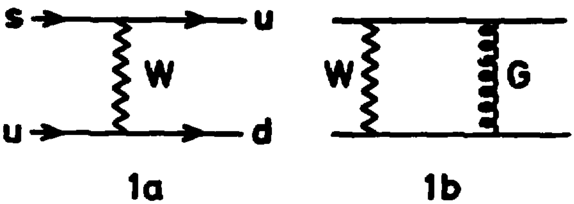


Fig.2