ON THE CONTRIBUTION OF NEW INTERACTIONS TO THE $\Delta I = 1/2$ RULE AND *CP* VIOLATION PARAMETERS

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Received 26 January 1987

New interactions which could appear as low-energy remnants of high-dimensional theories are studied in connection with their contribution to the $\Delta I = 1/2$ rule and *CP* violation parameters. The study is done in the framework of a particular superstring-inspired model; however, similar results could be obtained in a large class of models.

The attempt to unify all known interactions has recently led to speculations that such a unification might occur in higher dimensions. Very encouraging is the fact that superstring theories [1] can suggest even the gauge group in higher dimensions. It is particularly interesting to examine to which extent the new interactions that might occur, after dimensional reduction in the effective four-dimensional theory, could also help in understanding long-standing puzzles in low-energy physics, such as the $\Delta I = 1/2$ rule in hadronic weak decays [2]. Some time ago there was some hope that the $\Delta I = 1/2$ rule could be explained if the contribution of the penguin diagram in the $K \to 2\pi$ decay is taken into account [3]. However, this hope faded when it was realized, among other things, that the enhancement of the $\Delta I = 1/2$ amplitudes drives the ratio ϵ'/ϵ beyond the allowed experimental limits [4]. In this letter we examine the contribution to the $\Delta I = 1/2$ rule, coming from new interactions, taking into account the allowed bound on the *CP* violation parameter.

In most of the attempts to obtain a unified theory in four dimensions, after dimensionally reducing the high-dimensional theory, E_6 emerges as the gauge group [1]. This group is assumed to break down at some scale to $SU(3) \times SU(2) \times U(1)$, which eventually is hoped to break to $SU(3) \times U(1)_{em}$. Assigning the fermions in the 27-dimensional representation of E_6, one finds a new chiral D-superfield [5] which transforms as $(3, 1, -1/3)$ under $SU(3) \times SU(2) \times U(1)$. Allowing for some Z discrete symmetries to avoid several problems [5,6] (like the rapid proton decay), the allowed Yukawa couplings of the D-superfield with the usual quarks which appear in the superpotential are [5]

$$\lambda DQQ + \lambda' D^c u^c d^c + \text{h.c.}\,, \tag{1}$$

where Q is the left-handed quark doublet, u and d are the right-handed singlets and λ and λ' are the coupling constants.

Eq. (1) shows that D-scalars having the above couplings with the normal quarks, will contribute to the K^0–$\bar{K}^0$ mass difference through the box diagram (fig. 1a) and to the $K \to \pi\pi$ decay through the penguin diagram (fig. 2a). D-scalars seem to play the role of Higgs scalars in the standard model with nonminimal Higgs content. However, an important difference appears. D-scalars, being coloured, should not acquire a v.e.v. and therefore there is no reason for the couplings λ and λ' to be of $O(gm_f/\mu_W)$, where m_f is the mass of the ordinary quarks. In other words, the couplings λ and λ' could, in principle, be as large as the gauge coupling. In order to maximize any possible effect, due to the new couplings, we shall also assume that the couplings of the three generations of quarks couple with the same strength to the D's, which have the same mass, postulating in this way a super-D-GIM mecha-

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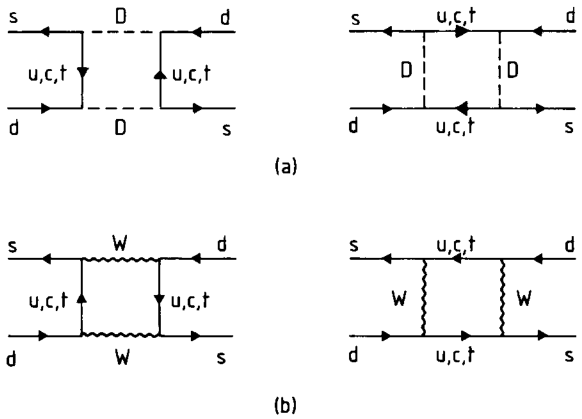

Fig. 1. Box diagram contributing to the $K^0-\bar{K}^0$ mass difference. (a) D-scalar exchange; (b) W exchange.

nism. This must also be obeyed by the right-handed D-scalars, appearing in the $D^c u^c d^c$ term. Without this assumption, λ has to be less than 9×10^{-2} from $K^0-\bar{K}^0$, less than 2×10^{-2} from $\mathrm{Im}(K^0-\bar{K}^0)$ and less than 3×10^{-3} from the neutron dipole moment [7]. A coupling of this order, as we shall see, gives a negligible enhancement in the $\Delta I = 1/2$ amplitude.

We start our analysis by evaluating the diagrams of figs. 2a and 2b. The effective hamiltonian H_D for the exchange of a D-scalar (fig. 2a), after using the equation of motion for the gluon field, takes the form

$$\mathcal{H}_D = \tfrac{1}{2} g_s \lambda^2 \frac{1}{m_G^2} \sum_{i=u,c,t} (U_{si} U^*_{di}) f_i \partial^\mu (\bar{s}_R \sigma_{\mu\nu} \lambda^a d_L)(\bar{u}\gamma^\nu \lambda^a u + \bar{d}\gamma^\nu \lambda^a d + \ldots)\,, \tag{2}$$

where g_s is the strong coupling constant, the U's are the matrix elements of the KM matrix

$$\begin{pmatrix} U_{ud} & U_{us} & U_{ub} \\ U_{cd} & U_{cs} & U_{cb} \\ U_{td} & U_{ts} & U_{tb} \end{pmatrix} = \begin{pmatrix} c_1 & s_1 c_3 & s_1 s_3 \\ -s_1 c_2 & c_1 c_2 c_3 - s_2 s_3 e^{i\delta} & c_1 c_2 c_3 + s_2 c_3 e^{i\delta} \\ -s_1 s_2 & c_1 s_2 c_3 + c_2 s_3 e^{i\delta} & c_1 s_2 s_3 - c_2 c_3 e^{i\delta} \end{pmatrix}\,, \tag{3}$$

and

$$f_i = \frac{m_i}{16\pi^2} \left\{ \frac{1}{M_D^2} \left(\frac{3}{2} + \ln \frac{m_i^2}{M_D^2} \right) \right\}\,, \tag{4}$$

where m_i is the mass of the ith (i = u, c, t) quark and M_D is the mass of the D-scalar. In eq. (4), it is assumed that $m_i^2 \ll M_D^2$. Finally, m_G in eq. (2) is an effective mass which appears in the gluon propagator. This must not sur-

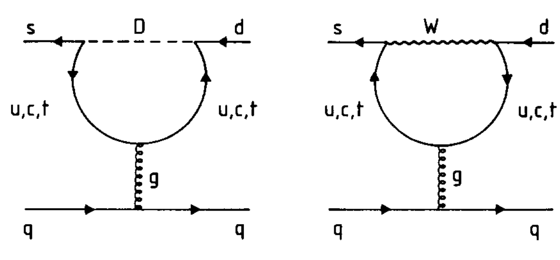

Fig. 2. Penguin diagrams contributing to $K \to 2\pi$ decays. (a) D-scalar exchange; (b) W exchange.

prise us, since, in contrast to the usual (W-exchange) penguin where the gluon pole cancels, here it is present. The same thing happens to the Higgs-exchange penguins [8].

Now the matrix elements of the four-quark operator are [9]

$$\langle\pi^+\pi^-|\partial^\mu(\bar{s}\sigma_{\mu\nu}(1-\gamma_5)\lambda^a d)(\bar{u}\gamma^\nu\lambda^a u+\bar{d}\gamma^\nu\lambda^a d+\ldots)|K^0\rangle=\frac{8}{3}\frac{f_\pi m_K^2 m_\pi^2}{m_d+m_u}\left(\frac{f_K}{f_\pi}+1+\frac{f_K}{f_\pi}\frac{m_K^2}{m_\sigma^2}\right), \tag{5}$$

where m_u, m_d are current algebra masses, f_π and f_K the pion and K-decay constants and m_K and m_σ the masses of the K and the 0^+ scalar meson ($m_\sigma = 700$ MeV).

For the W-exchange penguin, $\mathcal{H}_W$ is given by

$$\mathcal{H}_W=\frac{G_F}{\sqrt{2}}\frac{\alpha_s}{12\pi}\sum_{i=u,c,t}(U_{si}U_{di}^*)\ln\left(\frac{m_i^2}{\mu^2}\right)(\bar{s}\gamma_\mu(1-\gamma_5)\lambda^a d)(\bar{u}\gamma^\mu\lambda^a u+\bar{d}\gamma^\mu\lambda^a d+\ldots), \tag{6}$$

where μ is a typical hadronic mass. The matrix element of the four-quark operator is [8]

$$\langle\pi^+\pi^-|(\bar{s}\gamma^\mu(1-\gamma_5)\lambda^a d)(\bar{u}\gamma_\mu\lambda^a u+\bar{d}\gamma_\mu\lambda^a d+\ldots)=\frac{8}{3}\frac{f_K m_K^2 m_\pi^2}{m_d+m_u}\left(\frac{f_K}{f_\pi}+1+\frac{f_K}{f_\pi}\frac{m_K^2}{m_\sigma^2}\right). \tag{7}$$

Now we define the quantity F:

$$F=\mathrm{Re}\langle\pi^+\pi^-|\mathcal{H}_D|K^0\rangle/\mathrm{Re}\langle\pi^+\pi^-|\mathcal{H}_W|K^0\rangle, \tag{8}$$

which measures the relative contribution to the $K\to 2\pi$ amplitude due to D-scalar and W exchanges. Using eqs. (2), (5), (6) and (7) in eq. (8), we shall express the factor F as a function of δ (the angle in the KM matrix), λ and M_D. These parameters are constrained mainly from the experimental values of ϵ, ϵ'/ϵ and the $K^0-\bar{K}^0$ mass difference.

The ϵ parameter is defined by [10]

$$|\epsilon|=2^{-3/2}|\epsilon_m+2\xi|, \tag{9}$$

where ξ, in our case, is defined by

$$\xi=\frac{\mathrm{Im}\langle\pi^+\pi^-|\mathcal{H}_D|K^0\rangle+\mathrm{Im}\langle\pi^+\pi^-|\mathcal{H}_W|K^0\rangle}{\mathrm{Re}\langle\pi^+\pi^-|\mathcal{H}_D|K^0\rangle+\mathrm{Im}\langle\pi^+\pi^-|\mathcal{H}_W|K^0\rangle}, \tag{10}$$

and ϵ_m is given by

$$\epsilon_m=\mathrm{Im}\langle K^0|\mathcal{H}|\bar{K}^0\rangle/\mathrm{Re}\langle K^0|\mathcal{H}|\bar{K}^0\rangle. \tag{11}$$

The major contribution to $\mathrm{Im}\langle K^0|\mathcal{H}|\bar{K}^0\rangle$ and $\mathrm{Re}\langle K^0|\mathcal{H}|\bar{K}^0\rangle$ is given by the box diagram with a W-exchange (fig. 1b) and a D-exchange (fig. 1a). The hamiltonians $\mathcal{H}_W$ and $\mathcal{H}_D$ are given by

$$\mathcal{H}_W=-\frac{G_F}{\sqrt{2}}\frac{\alpha}{16\pi}\left(\frac{m_c}{M_W\sin\theta_W}\right)^2\Xi, \tag{12}$$

$$\mathcal{H}_D=-\frac{\lambda^4}{8M_D^4}\frac{m_c^2}{64\pi^2}\Xi, \tag{13}$$

where

$$\Xi=(U_{cs}U_{cd}^*)^2+(U_{ts}U_{td}^*)^2\frac{m_t^2}{m_c^2}+U_{cs}U_{cd}^*U_{ts}U_{td}^*\frac{2m_t^2}{m_t^2-m_c^2}\ln\frac{m_t^2}{m_c^2}. \tag{14}$$

Since the quark current is the same in both diagrams of fig. 1, the matrix elements of the current do not appear in ϵ_m.

Finally, we need ϵ'/ϵ, which is given by

$$|\epsilon'/\epsilon| = |A_2/A_0|\,|2\xi/(\epsilon_m + 2\xi)|\,, \tag{15}$$

where A_0, A_2 are the amplitudes for $\Delta I = 1/2$ and $\Delta I = 3/2$ transitions (experimentally, $A_2/A_0 \approx 1/20$).

We now proceed to our numerical analysis. As a simple test, we put $\lambda = 0$ and using typical values for the vacuum parameters ($s_1 = 0.23, s_2 = 0.06, s_3 = 0.04$) we fit the experimental value of ϵ [11]:

$$|\epsilon| = (2.28 \pm 0.05) \times 10^{-3}\,, \tag{16}$$

and find $|\epsilon'/\epsilon| \approx 5.4 \times 10^{-3}$, in agreement with a previous analysis [2] and within the experimental limits [11]

$$|\epsilon'/\epsilon| = (-4.6 \pm 5.3 \pm 2.4) \times 10^{-3}\,. \tag{17}$$

Next we include the D-scalar contribution, without QCD corrections for the purpose of the present analysis. In fig. 3 we present our results. In the $M_D - \delta$ plane we plot the contours of $|\epsilon|$ for the values 2.2×10^{-3} and 2.4×10^{-3}. The (almost) vertical line is a contour of $|\epsilon'/\epsilon|$ for the value 10×10^{-3}. The other vertical segments correspond to other contours of $|\epsilon'/\epsilon|$ for the indicated values. On the upper horizontal line of the graph, we have written the corresponding values of the factor F (being almost independent of δ, dependent only on M_D). In fig. 3a, we have $\lambda = 0.3$ and in fig. 3b we have $\lambda = 0.1$. In the first case ($\lambda = 0.3$), masses below ~80 GeV are excluded by the $K^0-\bar{K}^0$ mass difference (shaded region).

We see that for $\lambda = 0.3$ and the super-D-GIM mechanism in operation we have a sizeable enhancement of the $\Delta I = 1/2$ amplitude without upsetting the value of $|\epsilon'/\epsilon|$. However, for $\lambda = 0.1$, the enhancement is already negligible.

In conclusion, we have presented an analysis starting from the question to what extent new interactions emerging in superstring-inspired models could enhance the $\Delta I = 1/2$ amplitude in weak hadronic decays without violating the bounds of the *CP* violation parameters. We find that it is possible to have a sizeable enhancement within the allowed region of the parameters. Even though the coupling constant which gives an appreciable contribution to the $\Delta I = 1/2$ rule is rather large, it is interesting to note that, in contrast with the case where the Higgs boson is

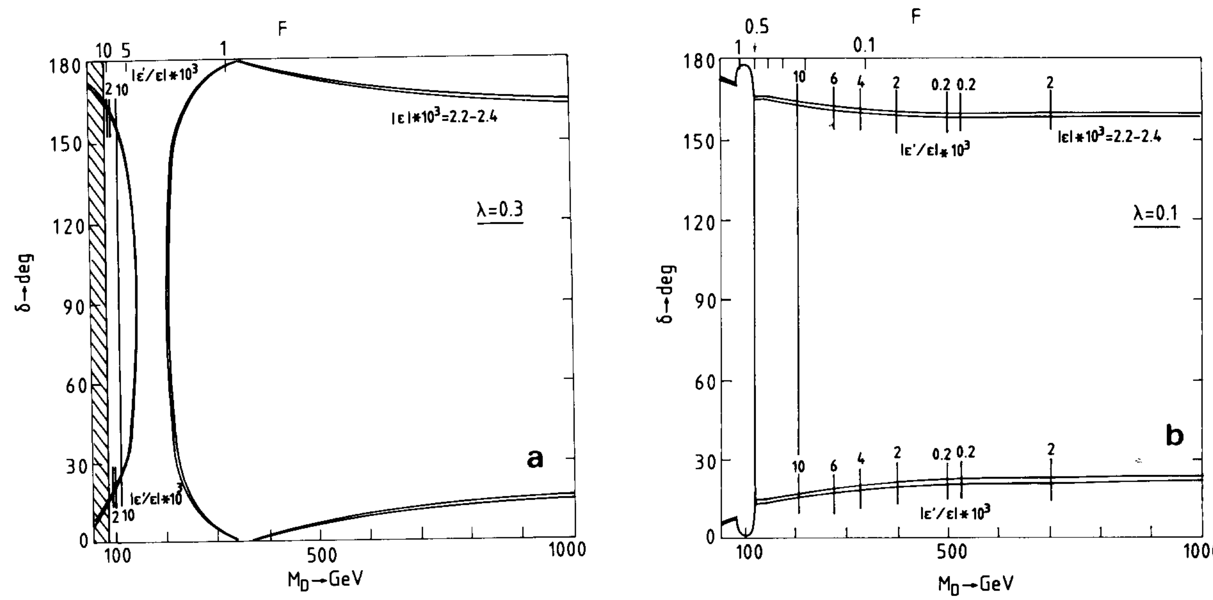

Fig. 3. Contours of $|\epsilon|$ and $|\epsilon'/\epsilon|$ in the $M_D - \delta$ plane. The enhancement factor F is shown on the upper horizontal axis of the graph. (a) $\lambda = 0.3$; (b) $\lambda = 0.1$.

used, such an enhancement can be achieved without giving a high $|\epsilon'/\epsilon|$ value. This observation could be used in similar models. Last but not least, we reemphasize the crucial role that our postulated super-D-GIM mechanism plays in our analysis.

We would like to thank J. Ellis for valuable discussions and for reading the manuscript. Two of us (E.P. and G.Z.) would like to thank CERN and the Greek Ministry of Research and Technology for their financial support.

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