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HIGGS INDUCED FLAVOUR CHANGING NEUTRAL INTERACTIONS
IN $SU(2)_L \times SU(2)_R \times U(1)$

G. Ecker

Institut für Theoretische Physik der
Universität Wien, Austria

W. Grimus

CERN - Geneva

and

H. Neufeld

Institut für Theoretische Physik der
Universität Wien, Austria

A B S T R A C T

Natural suppression of induced scalar contributions to $\langle \bar{K}_0 | M | K_0 \rangle$ in the simplest $SU(2)_L \times SU(2)_R \times U(1)$ gauge theory requires $M_H \gtrsim M_0 = 4-7$ TeV for the masses of neutral flavour changing Higgs particles. In addition, the quark mass matrices must be symmetric in a certain basis of weak eigenstates unless $M_H \gg M_0$.

Left-right symmetric gauge theories¹⁾ based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)$ are an attractive extension of the standard model²⁾. The observed structure of weak charged and neutral currents implies restrictions on the masses and mixings of heavy gauge vector bosons. Recently, Beall, Bander and Soni³⁾ have derived a lower bound of approximately 1.6 TeV for the heavy charged vector boson mass M_R from the $K_L K_S$ mass difference Δm_{LS} in a four quark model. In the meantime, it has been claimed that including the t quark and the Higgs bosons in the usual box-diagrams can bring that lower limit down considerably⁴⁾. However, in all those papers [except for a very recent preprint by Mohapatra, Senjanović and Tran⁵⁾ on which we shall comment at the end] the neutral Higgs exchange contribution to Δm_{LS} was not considered. We, therefore, want to investigate the neutral Higgs sector with as few assumptions as possible in order to derive the $|\Delta F| = 2$ flavour changing effective Hamiltonian induced by single Higgs exchange.

The Higgs sector consists of a number of multiplets

$$\phi_\alpha = \begin{pmatrix} \phi_{\alpha,11} & \phi_{\alpha,12} \\ \phi_{\alpha,21} & \phi_{\alpha,22} \end{pmatrix} \quad (\alpha = 1, \dots, n)$$

transforming as $(1/2, 1/2, 0)$ under the gauge group and of other scalar fields which do not couple to quarks. The neutral Hermitian field

$$\begin{aligned} \phi_0 &= \frac{i}{\sqrt{2} V} \sum_\alpha (V_\alpha^* \phi_{\alpha,11} - V_\alpha \phi_{\alpha,11}^\dagger + W_\alpha \phi_{\alpha,22}^\dagger - W_\alpha^* \phi_{\alpha,22}) = \\ &= \frac{\sqrt{2}}{V} \sum_\alpha (\text{Im } V_\alpha \text{Re } \phi_{\alpha,11} - \text{Re } V_\alpha \text{Im } \phi_{\alpha,11} - \\ &\quad - \text{Im } W_\alpha \text{Re } \phi_{\alpha,22} + \text{Re } W_\alpha \text{Im } \phi_{\alpha,22}) \end{aligned} \quad (1)$$

with

$$\langle \phi_\alpha \rangle_0 = \frac{1}{\sqrt{2}} \begin{pmatrix} V_\alpha & 0 \\ 0 & W_\alpha \end{pmatrix}, \quad V = \sqrt{\sum_\alpha (|V_\alpha|^2 + |W_\alpha|^2)} \simeq (\sqrt{2} G_F)^{-1/2}$$

contributes to the two neutral would-be-Goldstone bosons. Of the remaining $4n-1$ physical neutral Higgs fields of the general form

$$\begin{aligned} \phi_i &= \sum_\alpha (a_\alpha^i \text{Re } \phi_{\alpha,11} + b_\alpha^i \text{Im } \phi_{\alpha,11} + c_\alpha^i \text{Re } \phi_{\alpha,22} + d_\alpha^i \text{Im } \phi_{\alpha,22}) \\ i &= 1, \dots, 4n-1 \end{aligned} \quad (2)$$

we can choose $4n-2$ orthogonal fields with $\langle \phi_i \rangle_0 = 0$ (note that also $\langle \phi_0 \rangle_0 = 0$). This leaves us with a uniquely determined neutral field

$$\begin{aligned} \phi_1 &= \frac{\sqrt{2}}{V} \sum_{\alpha} (\text{Re } V_{\alpha} \text{Re } \phi_{\alpha,11} + \text{Im } V_{\alpha} \text{Im } \phi_{\alpha,11} + \text{Re } W_{\alpha} \text{Re } \phi_{\alpha,22} + \text{Im } W_{\alpha} \text{Im } \phi_{\alpha,22}) = \\ &= \frac{1}{\sqrt{2} V} \sum_{\alpha} (V_{\alpha}^* \phi_{\alpha,11} + V_{\alpha} \phi_{\alpha,11}^{\dagger} + W_{\alpha}^* \phi_{\alpha,22} + W_{\alpha} \phi_{\alpha,22}^{\dagger}) \end{aligned} \quad (3)$$

with non-vanishing VEV $\langle \phi_1 \rangle_0 = V$. Therefore, only $\langle \phi_1 \rangle_0$ contributes to the quark mass matrices and consequently the Yukawa couplings of the shifted field $\phi'_1 = \phi_1 - V$ are necessarily flavour diagonal:

$$-\mathcal{L}_Y^0 = \bar{u}_L \hat{M}_u u_R + \bar{d}_L \hat{M}_d d_R + \frac{\phi_1 - V}{V} (\bar{u}_L \hat{M}_u u_R + \bar{d}_L \hat{M}_d d_R) + h.c. + \quad (4)$$

+ Yukawa couplings of ϕ_i ($i = 2, 3, \dots, 4n-1$)

with diagonal, positive definite mass matrices $\hat{M}_u, \hat{M}_d$ and mass eigenfields u, d . The field ϕ'_1 is thus the analogue of the single Higgs field of the standard model whereas all other ϕ_i will in general have flavour changing Yukawa couplings.

Below the scale Λ_R of spontaneous parity violation there are $2n$ Higgs doublets under the effective gauge group $SU(2)_L \times U(1)$. Using an argument due to Georgi and Nanopoulos⁶⁾ we can always (unless additional symmetry requirements such as horizontal symmetries are interfering) arrange the Higgs potential such that all neutral Higgs fields except ϕ'_1 receive large masses. Actually, those masses cannot all be arbitrarily large because for masses $\gtrsim \Lambda_R$ the full symmetry $SU(2)_L \times SU(2)_R \times U(1)$ is operative. In order to suppress induced flavour changing neutral currents a realistic neutral Higgs sector will therefore contain only one low mass field with flavour diagonal Yukawa couplings whereas all other neutral Higgs have masses $\gtrsim \Lambda_R$.

Let us now turn to the minimal model⁷⁾ with $n = 1$ and in addition two scalar multiplets Δ_L, Δ_R transforming as $(1, 0, 1)$ and $(0, 1, 1)$, respectively. Studying the most general Higgs potential consistent with gauge and left-right symmetry one finds indeed a single Higgs field with naturally small mass, i.e.,

without assuming specific relations among the parameters of the potential. Up to mixing angles of the order of M_L/M_R which are model dependent this field is just ϕ'_1 . To the same accuracy there are two nearly degenerate heavy Higgs fields

$$\begin{aligned}\phi_2 &= \frac{1}{\sqrt{2}V} \left\{ \left| \frac{v}{w} \right| (w^* \phi_{22} + w \phi_{22}^\dagger) - \left| \frac{w}{v} \right| (v^* \phi_{11} + v \phi_{11}^\dagger) \right\} \\ \phi_3 &= \frac{i}{\sqrt{2}V} \left\{ \left| \frac{v}{w} \right| (w^* \phi_{22} - w \phi_{22}^\dagger) + \left| \frac{w}{v} \right| (v^* \phi_{11} - v \phi_{11}^\dagger) \right\}.\end{aligned}\quad (5)$$

The three remaining physical neutral fields are predominantly given in terms of the neutral components of Δ_L, Δ_R and have therefore negligible Yukawa couplings to the quarks. An analogous situation was found by Senjanović⁸⁾ in the original version of the model with doublet fields $X_{L,R}$ instead of triplets $\Delta_{L,R}$ in the limit $w = 0$.

Let us comment briefly on the limit $w = 0$. In the simplest model under consideration ($n=1$) the ratio $|w/v|$ is generically of the order m_b/m_t if the smallness of mixing angles is to have a natural explanation⁹⁾. Thus, $|w/v|$ is not necessarily very small and we prefer to keep $w \neq 0$ throughout this paper. Contrary to a frequently made claim the experimental or theoretical limits on the $W_L W_R$ mixing angle¹⁰⁾ do not force $|w| \ll |v|$ since this mixing angle is approximately given by $2(M_L/M_R)^2 |w/v|$.

The Yukawa couplings of the fields ϕ'_1, ϕ_2, ϕ_3 to quarks are given by

$$-\mathcal{L}'_Y = \frac{\phi'_1}{V} (\bar{u}_L \hat{M}_u u_R + \bar{d}_L \hat{M}_d d_R) + \sum_{i=2}^3 \phi_i (\bar{u}_L G_i^u u_R + \bar{d}_L G_i^d d_R) + h.c. \quad (6)$$

with

$$\begin{aligned}G_2^u &= -i G_3^u = (|v|^2 - |w|^2)^{-1} \left(-2 \frac{|vw|}{V} \hat{M}_u + V \frac{vw^*}{|vw|} K_L \hat{M}_d K_R^\dagger \right) \\ G_2^d &= i G_3^d = (|v|^2 - |w|^2)^{-1} \left(-2 \frac{|vw|}{V} \hat{M}_d + V \frac{v^* w}{|vw|} K_L^\dagger \hat{M}_u K_R \right)\end{aligned}\quad (7)$$

where $K_{L(R)} = U_{L(R)}^{u\dagger} U_{L(R)}^d$ are the left (right) handed mixing matrices and the unitary matrices $U_{L,R}^{u,d}$ diagonalize the quark mass matrices. The flavour mixing effective Hamiltonian at low energies induced by single Higgs exchange is therefore

$$\mathcal{H}_{\text{eff}}^{\Delta F \neq 0} = -\frac{1}{2} \sum_{i=2}^3 \frac{J_i(x)^2}{M_i^2} \quad (8)$$

where

$$J_i = \bar{u}_L G_i^u u_R + \bar{d}_L G_i^d d_R + \text{h.c.}$$

Let us concentrate on the $|\Delta S| = 2$ part

$$\begin{aligned} \mathcal{H}_{\text{eff}}^{|\Delta S|=2} = & -\frac{1}{2} \sum_{i=2}^3 M_i^{-2} \left\{ (G_{i,21}^d)^2 \bar{s}_L d_R \bar{s}_L d_R + (G_{i,21}^{d\dagger})^2 \bar{s}_R d_L \bar{s}_R d_L + \right. \\ & \left. + 2 G_{i,21}^d G_{i,21}^{d\dagger} \bar{s}_L d_R \bar{s}_R d_L \right\} + \text{h.c.} \end{aligned} \quad (9)$$

The real and imaginary parts of $\langle \bar{K}^0 | \mathcal{H}_{\text{eff}}^{|\Delta S|=2} | K^0 \rangle$ will contribute to Δm_{LS} and CP violating amplitudes, respectively. The more severe constraint applies to the imaginary part because experimentally $\text{Im} \langle \bar{K}^0 | M | K^0 \rangle / \text{Re} \langle \bar{K}^0 | M | K^0 \rangle \sim 10^{-3}$. Because of the relation

$$\langle \bar{K}^0 | \bar{s}_L d_R \bar{s}_L d_R | K^0 \rangle = \langle \bar{K}^0 | \bar{s}_R d_L \bar{s}_R d_L | K^0 \rangle \quad (10)$$

the only natural way to ensure a real $\langle \bar{K}^0 | \mathcal{H}_{\text{eff}}^{|\Delta S|=2} | K^0 \rangle$ is to require

$$G_{i,21}^{d\dagger} = G_{i,21}^{d*} \quad (11)$$

Assuming therefore symmetric Yukawa coupling matrices G_i^d one obtains

$$K_R = A K_L^* \quad (12)$$

where A is a diagonal unitary matrix. It is then straightforward to show that there exists a basis of weak eigenstates with symmetric quark mass matrices ^{*)}. The existence of such a basis is thus a natural condition for the absence of CP violation induced by $\mathcal{H}_{\text{eff}}^{|\Delta S|=2}$. Of course, such a condition will only be relevant if the Higgs masses M_i are not substantially larger than the lower bound we are going to derive from the contribution of $\mathcal{H}_{\text{eff}}^{|\Delta S|=2}$ to Δm_{LS} .

^{*)} -----
Consider as an example the basis transformation $q_L = U_L^d U_R^{dT} q'_L$, $q_R = q'_R$.

In the relevant case, where M_2, M_3 are nearly degenerate, the dominant contribution to $\Delta m_{LS} = m_{K_L} - m_{K_S} = -2\text{Re} \langle \bar{K}^0 | M | K^0 \rangle$ comes from

$$\begin{aligned} \mathcal{H}_{\text{eff}}^a &= - \sum_{i=2}^3 M_i^{-2} |G_{i,21}^d|^2 \bar{s}_L d_R \bar{s}_R d_L + \text{h.c.} = \\ &= - \frac{2 |G_{2,21}^d|^2}{M_H^2} \bar{s}_L d_R \bar{s}_R d_L + \text{h.c.} \end{aligned} \quad (13)$$

with $2M_H^{-2} = M_2^{-2} + M_3^{-2}$. We normalize the $\bar{K}^0 K^0$ matrix element of (13) to the value given by the vacuum insertion method¹¹⁾ by writing

$$\begin{aligned} \langle \bar{K}^0(p') | \bar{s}_L d_R \bar{s}_R d_L(x) | K^0(p) \rangle &= R \langle \bar{K}^0(p') | \bar{s}_L d_R \bar{s}_R d_L(x) | K^0(p) \rangle_{\text{vac}} = \\ &= - \frac{R F_K^2 m_K^2 e^{i(p'-p) \cdot x}}{2(2\pi)^3 \sqrt{4\omega_p \omega_{p'}}} \left(\frac{m_K^2}{(m_s + m_d)^2} + \frac{1}{6} \right) \end{aligned} \quad (14)$$

with

$$\langle \bar{K}^0(p) | \bar{s} \gamma^\mu \gamma_5 d(x) | 0 \rangle = \frac{i F_K p^\mu e^{ip \cdot x}}{\sqrt{2\omega_p} (2\pi)^3}.$$

The contribution of $\mathcal{H}_{\text{eff}}^a$ to Δm_{LS} is thus

$$\Delta m_{LS}(H_0)|_a = - \frac{R |G_{2,21}^d|^2}{M_H^2} F_K^2 m_K \left(\frac{m_K^2}{(m_s + m_d)^2} + \frac{1}{6} \right). \quad (15)$$

The remaining part of the effective Hamiltonian

$$\mathcal{H}_{\text{eff}}^b = - \frac{1}{2} \sum_{i=2}^3 M_i^{-2} \left\{ (G_{i,21}^d)^2 \bar{s}_L d_R \bar{s}_L d_R + (G_{i,21}^{d*})^2 \bar{s}_R d_L \bar{s}_R d_L \right\} + \text{h.c.} \quad (16)$$

does not contribute to Δm_{LS} in the limit $M_2 = M_3$. In general, one derives the bound

$$\left| \frac{\langle \bar{K}^0 | \mathcal{H}_{eff}^b | K^0 \rangle}{\langle \bar{K}^0 | \mathcal{H}_{eff}^a | K^0 \rangle} \right| < \frac{5}{6} \frac{|M_2^2 - M_3^2|}{M_2^2 + M_3^2} \quad (17)$$

where we have assumed the validity of the vacuum insertion method for the ratio of matrix elements. This bound is also valid in the more general case that the massive Higgs fields are arbitrarily orthogonal combinations of the fields ϕ_2, ϕ_3 in (5). Interpreting M_H as an average Higgs mass bounded by

$$11 M_2^{-2} + M_3^{-2} < 12 M_H^{-2} < M_2^{-2} + 11 M_3^{-2} \quad (18)$$

we can regard (15) as the total contribution $\Delta m_{LS}(H_0)$ due to $\mathcal{H}_{eff}^{|\Delta S|=2}$. Note that $\Delta m_{LS}(H_0)$ has a definite sign which is opposite to the experimental value if the sign of the relevant matrix element (14) is correctly given by the vacuum insertion method.

From (7) we obtain in the six-quark model

$$|G_{2,21}^d|^2 = \frac{V^2 f^2 m_c^2 s_1^2}{(|v|^2 - |w|^2)^2} \quad (19)$$

with

$$f m_c s_1 = \left| \sum_{i=u,c,t} m_i K_{L,1s}^* K_{R,id} \right|.$$

Taking into account (12), assuming small KM mixing angles¹²⁾ and neglecting m_u one finds approximately

$$f \simeq \left| 1 + s_2 (s_2 + s_3 e^{-i\delta}) e^{-i\alpha} \frac{m_t}{m_c} \right| \quad (20)$$

where α is an unknown phase. We, therefore, arrive at the lower bound

$$|\Delta m_{LS}(H_0)| > \left(\frac{m_c}{m_s + m_d} \right)^2 |R| f^2 s_1^2 \sqrt{2} G_F F_K^2 m_K^3 / M_H^2 \quad (21)$$

which is our main result.

In the original calculation of Gaillard and Lee¹¹⁾ in the standard model with four quarks, the experimental value $\Delta m_{LS,exp}$ was used to derive a value for m_c which was later confirmed by experiment. It has been a common attitude ever since to require that any additional contribution should at least not be bigger than $\Delta m_{LS,exp}$. If we therefore demand $|\Delta m_{LS}(H_0)| < \Delta m_{LS,exp} = 3.5 \times 10^{-15}$ GeV we get from (21)

$$M_H > 6.7 \sqrt{|R|} \text{ f TeV.} \quad (22)$$

With $f \approx 1$ and Ref. 13) $1/3 \lesssim |R| \lesssim 1$ one obtains

$$M_H \gtrsim 4-7 \text{ TeV.}$$

Note that this lower bound depends only on the ratio of current quark masses which is relatively well known¹⁴⁾. Since there are only two natural scales in a gauge theory based on $SU(2)_L \times SU(2)_R \times U(1)$ the lower bound (22) leads one to expect that the scale of parity violation is at least in the TeV range.

In a recent preprint⁵⁾, which we received after completion of this work, Mohapatra, Senjanović and Tran have also calculated the neutral Higgs contribution to Δm_{LS} . They argue that even for $M_H = 100$ GeV the huge Higgs contribution may be cancelled by equally large box-graph contributions which must be due to the t quark. While such a cancellation may be possible technically it requires a fine-tuning of at least three orders of magnitude between two terms which do not seem to be related in any obvious way. At least within the present framework of spontaneous symmetry breaking quark mass matrix elements and Higgs masses have different origins. Moreover, the estimate of m_c in Ref. 11) would be completely fortuitous because the t quark contribution must be three orders of magnitude bigger than $\Delta m_{LS,exp}$ in order to achieve the desired cancellation.

On the other hand, we have no completely convincing argument against such a cancellation, either. However, we disagree on the following two points. Our final result for $\Delta m_{LS}(H_0)$ differs from that of Ref. 5). In particular, $\Delta m_{LS}(H_0)$ always has a definite sign independent of the Higgs masses M_2, M_3 if the vacuum insertion method can at least be trusted for the ratio of matrix elements. Secondly, for $M_H \lesssim 1$ TeV and Hermitian quark mass matrices as assumed in Ref. 5) the effective Hamiltonian (9) will in general lead to much too large CP violation unless another fine-tuning of phases is envisaged. As already emphasized, the only natural way to exclude CP violation induced by (9) is via symmetric quark mass matrices. This can be achieved most simply by

requiring real Yukawa coupling matrices in the natural basis where these matrices are Hermitian due to left-right symmetry. In such models CP is violated only spontaneously.

Finally, let us briefly consider $|\Delta C| = 2$ transitions. The $D_L D_S$ mass difference Δm_D is bounded by Ref. 15)

$$|\Delta m_D| \leq \frac{1}{2} (\Gamma_S + \Gamma_L) \left(\frac{2\rho}{1-\rho} \right)^{1/2}. \quad (23)$$

Using experimental information¹⁶⁾ on the decay width $(\Gamma_S + \Gamma_L)/2$ and on the mixing parameter $\rho = \Gamma(D^0 \rightarrow \bar{D}^0 \rightarrow K^+ \pi^-) / \Gamma(D^0 \rightarrow K\pi)$ one gets approximately $|\Delta m_D| \lesssim 10^{-12}$ GeV to be compared to $\Delta m_{LS} = 3.5 \times 10^{-15}$ GeV. Moreover, the relevant couplings for $|\Delta C| = 2$ are smaller than for $|\Delta S| = 2$. Altogether, the resulting lower bound on M_H is reduced by approximately

$$\left(\frac{m_K^3 m_c^4 10^{-12} \text{ GeV}}{m_D^3 m_s^4 \Delta m_{LS}} \right)^{1/2} \sim 140$$

compared to the $|\Delta S| = 2$ case.

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REFERENCES

- 1) J.C. Pati and A. Salam, Phys. Rev. D10 (1975) 275;
R.N. Mohapatra and J.C. Pati, Phys. Rev. D11 (1975) 566 and 2558;
G. Senjanović and R.N. Mohapatra, Phys. Rev. D12 (1975) 1502.
- 2) S.L. Glashow, Nucl. Phys. 22 (1961) 579;
S. Weinberg, Phys. Rev. Lett. 19 (1967) 1264;
A. Salam, Proc. of the 8th Nobel Symposium, Stockholm, 1968, ed. N. Svartholm
(Almqvist and Wiksell, Stockholm, 1968), p. 367.
- 3) G. Beall, M. Bander and A. Soni, Phys. Rev. Lett. 48 (1982) 848;
G. Beall, "Constraints on a left-right symmetric electroweak theory",
Univ. of Cal. Irvine Thesis, (1982).
- 4) P.L. de Forcrand, "Hadronic weak processes and the mass of the right handed
boson", LBL-14692, (July 1982);
M. Denis, Louvain-la-Neuve preprint (1982);
U. Lercher, private communication.
- 5) R.N. Mohapatra, G. Senjanović and M.D. Tran, "Strangeness changing processes
and the limit on the right handed gauge boson mass", City College New York
preprint (1983).
- 6) H. Georgi and D.V. Nanopoulos, Phys. Lett. 82B (1979) 95.
- 7) R.N. Mohapatra and G. Senjanović, Phys. Rev. Lett. 40 (1980) 912 and
Phys. Rev. D23 (1981) 165.
- 8) G. Senjanović, Nucl. Phys. B135 (1979) 334.
- 9) G. Ecker, W. Grimus and W. Konetschny, Phys. Lett. 94B (1980) 381 and Nucl.
Phys. B177 (1981) 489.
- 10) CDHS Collaboration, Z. Phys. C12 (1982) 225;
I.I. Bigi and J.-M. Frère, Phys. Lett. 110B (1982) 255;
J.F. Donoghue and B.R. Holstein, Phys. Lett. 113B (1982) 382.
- 11) M.K. Gaillard and B.W. Lee, Phys. Rev. D10 (1974) 897.
- 12) M. Kobayashi and K. Maskawa, Progr. Theor. Phys. 49 (1973) 652.
- 13) J. Trampetić, " $\bar{K}^0 K^0$ matrix elements and constraints on the W_R mass",
MPI-PAE/PTh 43/82, (July 1982).
- 14) J. Gasser and H. Leutwyler, Phys. Rep. 87C (1982) 77.
- 15) R.L. Kingsley, S.B. Treiman, F. Wilczek and A. Zee, Phys. Rev. D11 (1975)
1919;
E.A. Paschos, Phys. Rev. Lett. 39 (1977) 858.
- 16) Review of Particle Properties, Phys. Lett. 111B (1982)