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Weinberg-Salam model and chiral symmetry of mesons

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Abstract. Weinberg-Salam model is considered in the light of SU(4) × SU(4) chiral symmetry for mesons. The Higgs doublet and the pseudoscalar mesons mix in this framework. Consequences of this mixing for the Higgs decays and the nonleptonic decays of the mesons are explored.

Keywords. Higgs field; mesons; Meson-Higgs potential; decays; Weinberg-Salam model; chiral symmetry; SU(4) × SU(4).

1. Introduction

It is well known that the low energy phenomenology of hadrons, especially the pseudoscalar mesons is successfully described by models or effective Lagrangians based on spontaneously-broken chiral groups which incorporate the results of current algebra and PCAC (Gasiorowicz and Geffen 1969; Lee 1972; Schechter and Singer 1975; Singer 1977). Now a gauge theory based on the group $SU_L(2) \times U(1)$ has emerged as the most likely candidate for a theory of electro-weak interactions (Weinberg 1967; Salam 1968). The hadronic currents in this model are actually currents corresponding to the chiral group. Hence, it is natural to consider effective Lagrangians based on chiral groups for hadrons in which the appropriate SU(2) × U(1) subgroup is gauged.

In this work, we consider an effective Lagrangian model for pseudoscalar mesons based on the group SU(4) × SU(4). The mesons transform nonlinearly as a (4, 4*) + (4*, 4) representation of the group. A subgroup $SU_L(2) \times U(1)$ of this is gauged (it is chosen such that we get the usual Cabibbo structure for the hadronic weak currents). We also have the usual Higgs doublet. The leptons transform as usual under $SU_L(2) \times U(1)$. The interactions involving the Higgs doublet and the mesons are invariant under $SU_L(2) \times U(1)$. It will be seen that the Higgs scalars and the mesons mix in this model. This has several interesting consequences. For example, the two body semileptonic decays of the mesons (K_{l_2}, π_{l_2} etc.) arise from the Yukawa couplings between the Higgs doublet and the leptons.

More interesting is the potential term in the Lagrangian involving both the Higgs scalars and the mesons. We consider only bilinear terms (by this we mean, linear in the matrix function of the pseudoscalars and the Higgs field). We see that the terms which explicitly break the chiral symmetry necessary to give masses to the pseudoscalars as well as mass differences among different isomultiplets are included in them. In other words, the spontaneous breaking of $SU_L(2) \times U(1)$ gauge group induces explicit strong symmetry-breaking, a fact which has been noticed earlier (Weinberg 1971). The decay rates of the Higgs particle can be explicitly calculated. The total decay rate for two meson final states is in the range 10^{17}–10^{18} sec^{-1} for Higgs mass between 10–100 GeV.

We also have strangeness and charm violating terms in the potential. These are over and above the 'current algebra' terms. There is no sign of $\Delta I = \frac{1}{2}$ dominance, but departures from it can be estimated. We put as an extra input, a nonleptonic interaction term which belongs to a $\underset{\sim}{20}$ dimensional representation of SU(4) (Altarelli *et al* 1975 a, b). We also take into account a pure $\Delta I = \frac{1}{2}$ term for *K*-decays (Shifman *et al* 1977). Presumably, they arise from QCD effects. We calculate the two-body nonleptonic decays of the pseudoscalars. We find that the decay rate for $K^+ \to \pi^+ \pi^0$ which arises from $\Delta I = 3/2$ effects is about 5·5 times the experimental number. It is known that the charm meson decays violate $\Delta I = \frac{1}{2}$ rule but even here we find that our model yields very large $\Delta I = 3/2$ contributions.

2. The model

The strong interaction part is an effective Lagrangian for the pseudoscalar mesons invariant under spontaneously broken $SU(4) \times SU(4)$. Define the meson-matrix:

$$M(\Phi) = \exp\left(\frac{2i}{f}\,\Phi\right), \tag{1}$$

Here

$$\Phi = \sum_{i=0}^{15} \frac{\lambda_i}{\sqrt{2}}\,\Phi_i$$

is the pseudoscalar meson-matrix. f is the (averaged) pseudoscalar meson decay constant. $M(\Phi)$ and $M^+(\Phi)$ transform according to the representations $(4, 4^*)$ and $(4^*, 4)$ respectively of $SU(4) \times SU(4)$ (Singer 1977).

The electroweak interactions are described by the $SU_L(2) \times U(1)$ gauge-invariant Weinberg-Salam model. For the mesons, the generators corresponding to the $SU_L(2)$ subgroup are

$$I_{1_L} + I_{13_L},\ I_{2_L} - I_{14_L} \text{ and } I_{3_L} + \frac{1}{\sqrt{3}} I_{8_L} - \frac{2}{\sqrt{6}} I_{15_L}$$

where I_{i_L} are the generators of $SU_L(4)$ group. The generator corresponding to the U(1) subgroup is

$$I_{3_R} + \frac{1}{\sqrt{3}} I_{8_R} - \frac{2}{\sqrt{6}} I_{15_R}$$

where I_{i_R} are the generators of SU_R (4) group. We have to rotate the meson fields

$$M \to \tilde{M} = u\,M.$$

so that we get the usual Cabibbo structure for the hadronic weak currents. We will not specify the rotation matrix further, right now. We have the complex Higgs doublet

$$\mathrm{H} = \begin{pmatrix} H^+ \\ H^0 \end{pmatrix},$$

and the leptons which transform in the usual fashion under $SU_L(2) \times U(1)$. Then our Lagrangian has the form:

$$\begin{aligned} \mathscr{L} = & \frac{f^2}{8} [\mathrm{Tr}\,\{D_\mu \tilde{M}\, D_\mu \tilde{M}^+\} + U\,(\det M + \det M^+)] \\ & + D_\mu \mathrm{H}^*\, D_\mu \mathrm{H} - V(H) - V(M, H) + \mathscr{L}_{\text{gaugefield}} \\ & + \mathscr{L}_{\text{lepton}} + \mathscr{L}_{\text{Yukawa}}. \end{aligned} \tag{2}$$

Here $D_\mu \tilde{M}$, the covariant derivative of $\tilde{M}$ is given by the expression:

$$D_\mu \tilde{M} = \partial_\mu \tilde{M} - i\,g\,\Lambda_i\, W_{\mu\,i}\, \tilde{M} + i\,g'\, B_\mu\, \tilde{M}\, \Lambda_3, \tag{3}$$

where $$\Lambda_1 = \tfrac{1}{2}\left[\begin{array}{cc|cc} 0 & 1 & & \\ 1 & 0 & \multicolumn{2}{c}{0} \\ \hline \multicolumn{2}{c|}{0} & 0 & 1 \\ & & 1 & 0 \end{array}\right], \quad \Lambda_2 = \tfrac{1}{2}\left[\begin{array}{cc|cc} 0 & -i & & \\ i & 0 & \multicolumn{2}{c}{0} \\ \hline \multicolumn{2}{c|}{0} & 0 & i \\ & & -i & 0 \end{array}\right]$$

$$\Lambda_3 = \tfrac{1}{2}\begin{bmatrix} 1 & & & \\ & -1 & & \\ & & -1 & \\ & & & +1 \end{bmatrix}$$

W_μ^i, B_μ are the gauge bosons corresponding to the SU_L (2) and U (1) groups and g, g' are the corresponding coupling constants. U is a strong interaction parameter. $V(M, H)$ is the part of the potential involving both H and M invariant under SU_L (2) $\times$ U (1). We will specify it later. As usual, the Higgs field has a nonvanishing vacuum expectation value

$$\langle \mathrm{H} \rangle_0 = \zeta = \begin{pmatrix} 0 \\ \xi \end{pmatrix}. \tag{5}$$

Then $W^\pm$ and Z (which is the same combination of W^3 and B as in W–S model), acquire masses given by the expressions:

$$M_W^2 = \frac{g^2\,\xi^2}{2} + \frac{g^2\,f^2}{4},$$

$$M_Z^2 = \frac{(g^2 + g'^2)}{g^2}\, M_W^2. \tag{6}$$

Note that the mass terms for $W^{\pm}$ and Z have an additional piece coming from the mesons. But the mass ratio M_W/M_Z remains the same.

In the Lagrangian of equation (2) we have mixing terms between the gauge bosons and the spin-0 fields given by the expression:

$$\mathscr{L}_{\text{mixing}} = -\frac{f}{4}[\text{Tr}\, \partial_\mu \Phi \, \{g\, \Lambda_i' \, W_\mu^i - g' \, \Lambda_3' \, B_\mu\}]$$

$$-\frac{i}{2}[\partial_\mu \mathrm{H}^{+} \, (g\, \tau_i \, W_\mu^i \zeta + g' \, B_\mu \, \zeta)] + \text{h.c.} \qquad (7)$$

Here, $\Lambda_i' = u^{+} \, \Lambda_i \, u.$ (8)

To get the Cabibbo structure of the hadronic weak currents, u should be chosen such that

$$\Lambda_-' = \begin{bmatrix} 0 & 0 & 0 & 0 \\ \cos\theta & 0 & 0 & -\sin\theta \\ \sin\theta & 0 & 0 & \cos\theta \\ 0 & 0 & 0 & 0 \end{bmatrix}, \; \Lambda_+' = \Lambda_-'^{\dagger}, \; \Lambda_3' = \Lambda_3, \qquad (9)$$

where $\Lambda_{\pm} = \Lambda_1 \pm i\, \Lambda_2$ and θ is the Cabibbo angle.

In the unitary gauge, the mixing terms are put equal to zero (in other words, the unphysical Goldstone bosons are eliminated). Then, we get the following relations:

$$H^{+} = -\frac{if}{2\xi}[\cos\theta\, (\Phi_{12} + \Phi_{43}) + \sin\theta\, (\Phi_{13} - \Phi_{42})], \qquad (10)$$

$$H^{-} = H^{+*},$$

$$H^{0*} - H^{0} = \frac{-if}{2\xi}[\Phi_{11} - \Phi_{22} - \Phi_{33} + \Phi_{44}].$$

In the ordinary $SU_L(2)\times U(1)$ theory, the right sides would have been zero. Here combinations of Higgs fields and the mesons are the Goldstone bosons. This is what we mean when we say that Higgs fields and the mesons mix in this model.

3. Semileptonic decays

Consider any one generation of leptons *e.g.* electron and its neutrino. $L = \begin{pmatrix} \nu_e \\ e_L \end{pmatrix}$ is an isodoublet and e_R is an isosinglet. Now, in the Weinberg-Salam model there are no direct couplings between the hadrons and the leptons. Our model is in the spirit of an effective lagrangian and we do not consider direct couplings of the mesons with leptons. The Yukawa interaction between the leptons and the Higgs field is given by the expression:

$$\mathscr{L}_Y = -\, G_\phi^e \, (\bar{L}\, \mathrm{H}\, e_R + \bar{e}_R \, \mathrm{H}^{\dagger} \, L). \qquad (11)$$

Defining the shifted field H' by

$$\mathbf{H} = \mathbf{H}' + \xi \tag{12}$$

$$\mathscr{L}_Y = - G^e_\phi \, \xi \, (\bar{e}_L e_R + \bar{e}_R e_L) - G^e_\phi (\bar{e}_L H^{0'} e_R + \text{h.c.}) - G^e_\phi (\bar{\nu}_L e_R H^+ + \text{h.c.}). \tag{13}$$

When we use the expression for H^+ and $H^{0*} - H^{0*}$ in terms of the mesons given by (10), the second and third terms will induce Yukawa couplings between the mesons and the leptons, which we call $\mathscr{L}_{YM}$. It is given by the expression

$$\mathscr{L}_{\text{YM}} = i\sqrt{2}\, m_e G \bar{\nu}_L e_R \{f\cos\theta\,(\Phi_{12}+\Phi_{43}) + f\sin\theta\,(\Phi_{13}-\Phi_{42})\} + \text{h.c.} + i\frac{m_e Gf}{\sqrt{2}}(\bar{e}_R e_L - \bar{e}_L e_R)\{\Phi_{11} - \Phi_{22} - \Phi_{33} + \Phi_{44}\}. \tag{14}$$

Here we have used the relation,

$$\xi^2 = \frac{G^{-1}}{2\sqrt{2}} - \frac{f^2}{2} \approx \frac{G^{-1}}{2\sqrt{2}}, \tag{15}$$

This expression is the same as in the current-current Lagrangian (including neutral current couplings) when we use PCAC and Dirac equation for the electrons (Marshak *et al* 1969). In effect, the unitary gauge conditions take the place of PCAC.

Semileptonic decays involving more than one meson proceed as usual through the W-exchange. Such decays have been worked out in the framework of chiral dynamics (Aubrecht and Slanec 1981).

4. Meson-Higgs potential

To get the structure of the Meson-Higgs potential for an arbitrary orientation $\tilde{M}$ (consistent with the Cabibbo picture), it is simpler to deal with quarks and then go over to the physical meson fields.

Let q_1, q_2, q_3, q_4 denote the up, down, strange and charmed quarks respectively. Under $SU_L(2) \times U(1)$,

$$Q_{1_L} = \begin{pmatrix} q_1 \\ q_2 \end{pmatrix}_L \text{ and } Q_{2_L} = \begin{pmatrix} q_3 \\ q_4 \end{pmatrix}_L$$

are $SU_L(2)$ doublets. $q_{1_R}, q_{2_R}, q_{3_R}$ and q_{4_R} are singlets. The most general form for the Yukawa interaction $V_0(Q, H)$ (invariant under $SU_L(2) \times U(1)$) is:

$$V_0(Q, H) = \mathbf{H}^\dagger [(\bar{q}_{2_R}(\beta_1 Q_{1_L} + \beta_2 Q_{2_L}) + q_{3_R}(\beta_3 Q_{1_L} + \beta_4 Q_{2_L}) + \tilde{\mathbf{H}}^\dagger \bar{q}_{1_R}(\beta_5 Q_{1_L} + \beta_6 Q_{2_L}) + \bar{q}_{4_R}(\beta_7 Q_{1_L} + \beta_8 Q_{2_L})] + \text{h.c.} \tag{16}$$

where $\tilde{\mathbf{H}} = i\,\sigma_2 \mathbf{H}^*$.

As the Higgs field has a nonzero vacuum expectation value, this includes quark mass terms. It is necessary to diagonalize the quark mass matrix by rotating the quark fields. After diagonalization,

$$V_0(Q,H)\rightarrow V(Q,H)=\frac{(H^{0*}+H^0)}{2\xi}[B_1\bar{q}_1q_1+B_2\bar{q}_2q_2+B_3\bar{q}_3q_3+B_4\bar{q}_4q_4]$$

$$+\frac{(H^{0*}+H^0)}{2\xi}[B_1\bar{q}_1\gamma_5q_1-B_2\bar{q}_2\gamma_5q_2-B_3\bar{q}_3\gamma_5q_3+B_4\bar{q}_4\gamma_5q_4]$$

$$+\Big\{\frac{H^-}{\xi}[B_2\cos\theta\,\bar{q}_{2_R}q_{1_L}-B_1\cos\theta\,\bar{q}_{2_L}q_{1_R}$$

$$-B_2\sin\theta\,\bar{q}_{2_R}q_{4_L}+B_4\sin\theta\,\bar{q}_{2_L}q_{4_R}$$

$$+B_3\sin\theta\,\bar{q}_{3_R}q_{1_L}-B_1\sin\theta\,\bar{q}_{3_L}q_{1_R}$$

$$+B_3\cos\theta\,\bar{q}_{3_R}q_{4_L}-B_4\cos\theta\,\bar{q}_{3_L}q_{4_R}]+\text{h.c.}\Big\}. \tag{17}$$

Here B_i are functions of β_i and θ is the Cabibbo angle (which is the relative rotation of (q_2, q_3) and (q_1, q_4).

To get the potential in terms of the mesons, we have to make the substitutions,

$$\bar{q}_{i_R}q_{j_L}\rightarrow M_{ji},\quad \bar{q}_{i_L}q_{j_R}\rightarrow M^+_{ji}.$$

Using the Goldstone boson conditions expressed by (10) and defining the physical Higgs field χ by,

$$\chi=\frac{H^{0\prime}+H^{0\prime *}}{\sqrt{2}}=\frac{H^0+H^{0\prime *}}{\sqrt{2}}-\frac{2\xi}{\sqrt{2}},$$

$$V(Q,H)\rightarrow -V(M,H)$$

$$=\frac{f^2}{8}\sum_{i=1}^{4}\{A_i(M+M^+)_{ii}\}+\frac{f^2}{8}\frac{\chi}{\xi\sqrt{2}}\sum_{i=1}^{4}\Big\{\{A_i(M+M^+)_{ii}\}$$

$$+\Big[\frac{if^3}{16\xi^2}\{\cos\theta\,(\Phi_{21}+\Phi_{34})+\sin\theta\,(\Phi_{31}-\Phi_{24})\}$$

$$\times\{A_2\cos\theta\,M_{12}-A_1\cos\theta\,M^+_{12}-A_2\sin\theta\,M_{42}$$

$$+A_4\sin\theta\,M^+_{42}+A_3\sin\theta\,M_{13}-A_1\sin\theta\,M^+_{13}$$

$$+A_3\cos\theta\,M_{43}-A_4\cos\theta\,M^+_{43}\}+\text{h.c.}\Big]$$

$$+\frac{if^3}{32\,\xi^2}(\Phi_{11}-\Phi_{22}-\Phi_{33}+\Phi_{44})\,\{A_1\,(M-M^+)_{11}$$

$$-A_2\,(M-M^+)_{22}-A_3\,(M-M^+)_{33}+A_4\,(M-M^+)_{44}\} \tag{18}$$

Here we have defined A_i by

$$A_i=\frac{8}{f^2}B_i,$$

for later convenience. The first term in $V(M, H)$ explicitly breaks the SU(4) $\times$ SU(4) symmetry of the mesons. A_i are related to the meson masses (Singer 1977). The second term describes the coupling of the physical Higgs field to mesons. It is used to calculate the Higgs decays in the next section. The third term includes strangeness and charm-violating parts and contributes to the nonleptonic decays of mesons treated in the sixth section. Both the third and fourth terms give negligible corrections to the meson masses.

5. Higgs decays

When we expand the second term in $V(M, H)$ in (18), we get the following effective Lagrangian describing the coupling of the Higgs field to the mesons:

$$\mathscr{L}(\chi)=\frac{-\chi}{\sqrt{2}\xi}[m_\pi^2\,\pi^+\pi^-+m_K^2(K^+K^-+K^0\bar{K}^0)$$

$$+m_D^2\,(D^+D^-+D^0\,\bar{D}^{-0})+m_F^2F^+F^-+2\,A_1\,\Phi_{11}^2+2\,A_2\,\Phi_{22}^2$$

$$+2\,A_3\,\Phi_{33}^2+2\,A_4\,\Phi_{44}^2]+\ldots \tag{19}$$

For the charged mesons, the coupling is proportional to the square of the mass. For the diagonal fields, π^0, η, η', η'', it depends on the details of the mixing. The previous calculations of Higgs decays relied on broken scale invariance arguments, treating Higgs as a dilaton (Ellis *et al* 1976). Here the decay rates involve only known parameters and can be readily computed. We have considered only the two-body decays. The results are summarised in table 1 and figure 1. The total decay rate in the two-

Table 1. Higgs decay rates for two meson final states.

Higgs mass (GeV).	Decay rate for charged final states (sec^{-1})	Total decay rate (sec^{-1})
0·5	$2{\cdot}95\times10^{14}$	$4{\cdot}40\times10^{14}$
1	$5{\cdot}32\times10^{15}$	$10{\cdot}55\times10^{15}$
5	$17{\cdot}67\times10^{17}$	$35{\cdot}45\times10^{17}$
10	$12{\cdot}68\times10^{17}$	$35{\cdot}8\times10^{17}$
20	$6{\cdot}73\times10^{17}$	$19{\cdot}59\times10^{17}$
100	$1{\cdot}37\times10^{17}$	$4{\cdot}02\times10^{17}$

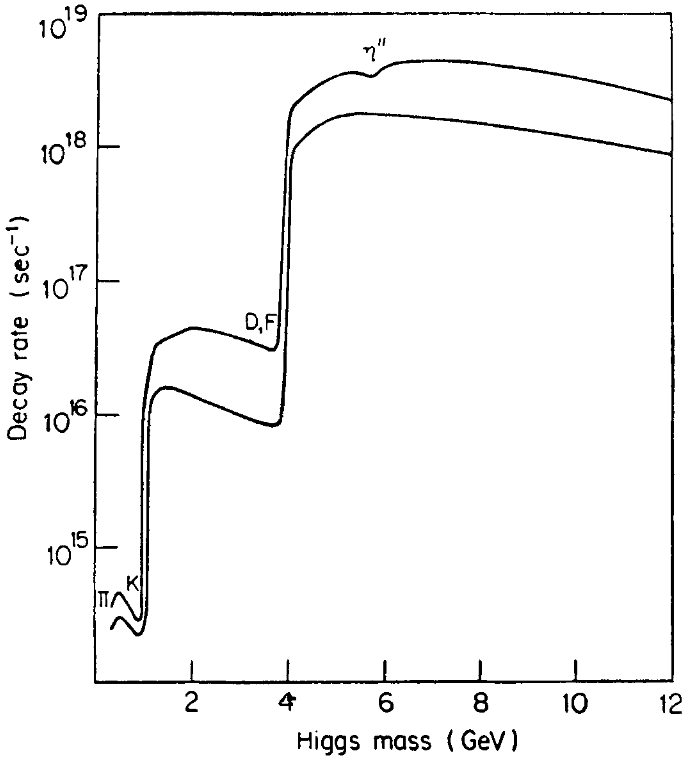

Figure 1. Higgs decay rate as a function of its mass. The decay rate is on a logarthmic scale. Various thresholds are indicated.

body meson channel is in the range 10^{17}–10^{18} sec^{-1} for Higgs mass between 10 and 100 GeV.

Due to the mixing between the Higgs field and the mesons, the Higgs potential $V(H)$ in (2) also contributes to the Higgs decay. But it can be verified that this contribution is substantial only for very high Higgs mass $\sim 10^6$ GeV where the ordinary tree approximation itself is not valid.

6. Nonleptonic decays

The third term in $V(M, H)$ in (18) contains strangeness and charm-violationg parts and contribute to the nonleptonic decays. It can be readily verified that it does not exhibit the experimentally observed $\Delta I = \frac{1}{2}$ rule or the 'octet dominance' in strange particle decays. For long, this rule has been a puzzling feature of nonleptonic decays as the effective nonleptonic Hamiltonian contain terms which transform as $\underline{8}$ and $\underline{27}$ representations of SU(3) with comparable strengths (see Marshak *et al* 1969 for earlier work). It had been pointed out by Wilson (1969) that short-distance effects due to strong interactions may enhance the octet part of the nonleptonic Hamiltonian in comparison with the $\underline{27}$ part. In an asymptotically-free theory of strong interactions like QCD, these effects can be explicitly calculated. Such a calculation has been performed and it has been shown that the octet part is indeed enhanced (Gaillard and Lee 1974; Altarelli and Maiani 1974). It has also been shown that when the charm quantum number is included, the enhanced part corresponds to a $\underline{20}$ dimensional representation of SU(4) (Altarelli *et al* 1975). But the magnitude of enhancement is much smaller than what one requires. It was pointed out later that there are additional contributions due to what are known as 'penguin' diagrams which are $(V-A)\otimes(V+A)$ in nature and have pure $\Delta I=\frac{1}{2}$ structure (Shifman *et al* 1977). Though it cannot be claimed that the strong interaction effects are completely understood two things stand out: (i) it is very likely that the 'penguin' type contributions dominate the strange particle decays, (ii) they play an insignificant role in the decays of charmed particles (see for example Sanda 1980). We will assume in the following that the $\Delta I=\frac{1}{2}$ effects are incorporated in an effective weak Lagrangian density, $\mathscr{L}_W$ proportional to a $\underline{20}$ dimensional representation and an additional $\Delta I=\frac{1}{2}$ term arising from the 'penguins' contributing to the strange particle decays alone. We will will not specify the latter further. Following the notation of Singer (1977),

$$\mathscr{L}_W=\frac{G}{2\sqrt{2}}\bar{X}\,[\sin\theta\cos\theta\,(j_{2a_\alpha}j_{a3_\alpha}-2\,j_{24_\alpha}j_{43_\alpha}$$

$$+2\,j_{23_\alpha}j_{44_\alpha})+\cos^2\theta\,(j_{43_\alpha}j_{21_\alpha}-j_{23_\alpha}j_{41_\alpha})$$

$$-\sin^2\theta\,(j_{42_\alpha}j_{31_\alpha}-j_{41_\alpha}j_{32_\alpha})+\sin\theta\cos\theta\,(j_{43_\alpha}j_{31_\alpha}-j_{41_\alpha}j_{33_\alpha}$$

$$-j_{42_\alpha}j_{21_\alpha}+j_{41_\alpha}j_{22_\alpha})]+\text{h.c.} \tag{20}$$

Here X is a dimensional parameter to be fixed from the rate for the decay $D^0\to c^-\pi^+$. The currents j_{ba_α} are the left-handed currents;

$$j_{ab_\alpha}=(V_{ab})_\alpha+(P_{ab})_\alpha-\tfrac{1}{4}\,\delta_{ab}\,(V_{cc_\alpha}+P_{cc_\alpha}), \tag{21}$$

where the vector currents V_{ab_α} and the axial vector currents P_{ab_α} are given by the expressions:

$$(V_{ab})_\alpha=\frac{if^2}{4}\,[M,\partial_\alpha M^+]_{ab}, \tag{22}$$

$$(P_{ab})_\alpha=\frac{if^2}{4}\,\{M,\partial_\alpha M^+\}_{ab}. \tag{23}$$

So, our nonleptonic Lagrangian is:

$$\mathscr{L}_{NL}=\mathscr{L}_{\text{penguin}}+\mathscr{L}_W+\frac{if^3}{16\xi^2}\,[\{\cos\theta\,(\Phi_{21}+\Phi_{34})+\sin\theta\,(\Phi_{31}-\Phi_{24})\}$$

$$\times\{A_2\cos\theta\,M_{12}-A_1\cos\theta\,M_{12}^+-A_2\sin\theta\,M_{42}+A_4\sin\theta\,M_{42}^+$$

$$+A_3\cos\theta\,M_{13}-A_1\sin\theta\,M_{13}^++A_3\cos\theta\,M_{43}-A_4\cos\theta\,M_{43}^+\}+\text{h.c.}] \tag{24}$$

Here we confine ourselves to the two-body nonleptonic decays of mesons. It is straightforward to calculate the decay rates and they are tabulated in table 2. In general, the model gives rise to substantial violation of the $\Delta I = \frac{1}{2}$ rule. The predicted value for the decay rate for $K^+ \to \pi^+ \pi^0$ $(9.13 \times 10^7 \text{ sec}^{-1})$ is about 5.5 times the experimental number $(1.629 \times 10^7 \text{ sec}^{-1})$. Also, the decay rate for $K_S^0 \to \pi^0 \pi^0$ $(0.77 \times 10^{10}$

Table 2. The decay rates for two body nonleptonic decays of the mesons.

Process	Decay rate (10^{10} sec^{-1})	Process	Decay rate (10^{10} sec^{-1})
$K_s^0 \to \pi^+ \pi^-$	1·10 (input)		
$K_s^0 \to \pi^0 \pi^0$	0·77		
$K^+ \to \pi^+ \pi^0$	$9·13 \times 10^{-3}$		
D^0 decays		*D^+ decays*	
Mode		Mode	
$K^- \pi^+$	5 (input)	$\bar{K}^0 \pi^+$	19·9
$K^+ \pi^-$	0·02	$K^+ \pi^0$	0·37
$K^+ K^-$	0·29	$K^+ \eta$	$2·8 \times 10^{-3}$
$\pi^+ \pi^-$	0·39	$K^+ \eta'$	$1·4 \times 10^{-3}$
$\bar{K}^0 \pi^0$	2·5	$K^0 \pi^+$	1·15
$\bar{K}^0 \eta$	1·03	$K^+ \bar{K}^0$	0·29
$\bar{K}^0 \eta'$	0·23	$\pi^+ \pi^0$	0·65
$K^0 \pi^0$	6×10^{-3}	$\pi^+ \eta$	0·86
$K^0 \eta$	$4·2 \times 10^{-3}$	$\pi^+ \eta'$	0·02
$K^0 \eta'$	$2·4 \times 10^{-3}$		
$\pi^0 \pi^0$	0·14	Total	23·24
$\eta \eta$	0·07		
$\pi^0 \eta$	0·06		
$\pi^0 \eta'$	0·14		
$\eta \eta'$	0·02		
Total	9·89		
F^+ decays		*η'' decays*	
Mode		Mode	
$\pi^+ \eta$	2·08	$K^- \pi^+$	$7·6 \times 10^{-4}$
$\pi^{\dagger} \eta'$	3·15	$K^0 \pi^0$	$3·8 \times 10^{-4}$
$K^+ K^0$	3·95	$K^0 \eta$	5×10^{-5}
$K^+ K^0$	0·06	$K^0 \eta'$	1×10^{-3}
$K^+ \pi^0$	0·10	$F^+ \pi^-$	18·29
$K^+ \eta$	0·78	$D^- K^+$	$3·9 \times 10^{-3}$
$K^+ \eta'$	0·31	$D^0 K^0$	$8·4 \times 10^{-3}$
$K^0 \pi^+$	12·71	$F^- K^+$	0·05
		$D^0 \pi^0$	0·17
Total	23·15	$D^0 \eta$	0·19
		$D^0 \eta'$	—
		$D^- \pi^+$	0·08
		$F^- \pi^+$	7·1
		Total	25·88

sec^{-1}) is noticeably higher than the experimental value (0.35×10^{10} sec^{-1}). Our results for charmed meson decays are at variance with the available experimental data (Particle Data Group 1982.) For the ratio of the decay rates of these particles, the data assign the values

$$\frac{\Gamma(D^+)}{\Gamma(D^0)} \approx 0.53, \frac{\Gamma(F^+)}{\Gamma(D^0)} \approx 2.18,$$

whereas our calculations indicate;

$$\frac{\Gamma(D^+)}{\Gamma(D^0)} \approx 2.35, \frac{\Gamma(F^+)}{\Gamma(D^0)} \approx 2.34.$$

We also get far too high a value for the decay rate for $D^+ \to \bar{K}^0 \pi^+$ (19.9×10^{10} sec^{-1}) compared with the experimental number (2×10^{10} sec^{-1}) (If the $\Delta I = \frac{1}{2}$ rule is applicable to the charm decays, this decay is forbidden!) All these discrepancies can be ascribed to the large $\Delta I = 3/2$ terms in our Lagrangian.

7. Conclusions

We have attempted to construct an effective Lagrangian for mesons in the framework of Weinberg-Salam model and chiral symmetry of the strong interactions. The consequences of mixing of Higgs fields and meson fields in such an approach were investigated. Some simple assumptions were made regarding the meson-Higgs potential. It was found that we get extra contributions for nonleptonic decays over and above the 'current-algebra' terms. Unfortunately these include large $\Delta I = 3/2$ terms. A more general method to construct the effective Lagrangian is called for.

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