ISTITUTO NAZIONALE DI FISICA NUCLEARE

Sezione di Trieste

INFN/AE-89/11

18 ottobre 1989

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LOOKING FOR A LIGHT HIGGS BOSON IN Φ AND Ψ RADIATIVE DECAYS

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Abstract

We discuss the possibility of producing a light Higgs boson in the ϕ radiative decay, using a vector meson dominance approach. The result we find for the branching ratio of this process should be of interest to a ϕ factory. We also extend our considerations to the analogous ψ radiative decay.

Considerable attention has been given recently to the question of the existence of a very light Higgs boson[1,2]. Such a possibility occurs in most extensions of the standard model[3], and also in the minimal model itself for a selected range of top quark masses[4]. A variety of processes have been considered as being particularly suitable to search for signatures of a low mass Higgs boson, and the resulting experimental limits have been critically surveyed. These include some nuclear transitions, kaon two- and three-body decays, and beauty decays. In particular the analysis of [2] has suggested that the regions 14 MeV $< m_H <$ 200 MeV and $m_H >$ 700 MeV could still be accessible to the light Higgs boson mass. Indeed the derivation of experimental constraints from the data is not completely unambiguous, as it requires model independent theoretical determinations of the Higgs production and decay matrix elements, which are not available in general. Therefore in this situation it is important to examine as many different processes as possible.

In this note we would like to point out that a relevant process in this context might be represented by the two-body radiative decay $\phi(1.02) \to \gamma H$, if this could be studied with enough sensitivity. In what follows we present a very simple estimate of this decay in the minimal version of the standard model, making use of the familiar vector meson dominance framework. We also extend our calculation to the $J/\psi(3.1)$ in order to obtain a qualitative estimate for the analogous transition $J/\psi \to \gamma H$.

We define the transition amplitude for the radiative process as

$$A(\phi(p,\eta) \to \gamma(q,\epsilon)H(k)) = \epsilon_\nu(q)\eta_\mu(p)\left[g_{\mu\nu} - \frac{q_\mu P_\nu}{(qP)}\right] G(k^2), \tag{1}$$

where $P = p + q$, $k = p - q$, and $k^2 = m_H^2$. Correspondingly the decay width is given by:

$$\Gamma(\phi \to \gamma H) = \frac{1}{24\pi m_\phi}\left(1 - \frac{m_H^2}{m_\phi^2}\right)|G|^2. \tag{2}$$

An expression of the amplitude, which is particularly convenient to the subsequent application of vector meson dominance of the photon, is obtained by reducing the latter from the final state. This gives:

$$A = e\left(\sqrt{2}G_F\right)^{1/2}\epsilon_\nu(q)\int d^4x \exp(iqx) < 0|T\left(j_\nu^{em}(x), \sum_i m_i\bar{q}_i q_i\right)|\phi(p,\eta)>, \tag{3}$$

where j_ν^{em} is the hadronic electromagnetic current, the q_i are the quark fields and the sum runs over all quark flavors. In eq.(3) we retain in the interaction hamiltonian density only the Higgs coupling to quarks via the quark mass term, which is dominant in the present case of no flavour change. To separate out the contribution from heavy quarks we write[5,6]:

$$\sum_i m_i\bar{q}_i q_i = -\frac{\beta - \beta_L}{\beta_L}\,\theta + \left(1 + \frac{\beta - \beta_L}{\beta_L}\right)\sum_L m_l\bar{q}_l q_l, \tag{4}$$

where θ is the trace of the energy-momentum tensor $\theta_{\alpha\beta}$, and L is for light quarks ($L =$ u, d, s in the present case), so that

$$\beta = -\frac{33 - 2n_f}{24\pi}\alpha_S = \beta_L + \frac{2n_{heavy}}{24\pi}\alpha_S, \tag{5}$$

with

$$\beta_L = -\frac{33 - 2n_L}{24\pi}\alpha_S. \tag{6}$$

Thus the amplitude G of interest in eq.(1) is split according to eq.(4) into two components that we must estimate separately.

For the contribution of the θ-term we start from:

$$\epsilon_\nu(q)T_{\nu\alpha\beta}(p,q,k) = \epsilon_\nu(q)\int d^4x \exp(iqx) < 0|\mathrm{T}\Big(j_\nu^{em}(x), \theta_{\alpha\beta}(0)\Big)|\phi(p,\eta)>, \tag{7}$$

and expand the tensor $T_{\nu\alpha\beta}$ into invariant amplitudes subject to the conditions

$$k_\alpha T_{\nu\alpha\beta} = k_\alpha T_{\nu\beta\alpha} = 0. \tag{8}$$

The result is:

$$\begin{aligned} T_{\nu\alpha\beta} =& \eta_\nu[A_1 l_\alpha l_\beta + A_2(k_\alpha k_\beta - k^2 g_{\alpha\beta})] \\ &+ A_3[l_\alpha(\eta_\beta k_\nu - (k\eta)g_{\beta\nu}) + (\alpha \leftrightarrow \beta)] \\ &+ A_4[\eta_\alpha(k_\beta k_\nu - k^2 g_{\beta\nu}) + (k\eta)(g_{\beta\nu}k_\alpha - g_{\alpha\beta}k_\nu) + (\alpha \leftrightarrow \beta)], \end{aligned} \tag{9}$$

where we have defined $l_\alpha = P_\alpha - \frac{(kP)}{k^2}k_\alpha$ orthogonal to k_α, and $A_i = A_i(p^2, q^2, k^2)$. By further imposing current conservation:

$$q_\nu T_{\nu\alpha\beta} = 0, \tag{10}$$

the tensor $T_{\nu\alpha\beta}$ can be written in terms of only one invariant amplitude, which we choose to be A_1. The final result at $q^2 = 0$ is, after taking the trace:

$$\epsilon_\nu T_{\nu\alpha\alpha} = A_1(\mathrm{m}_\phi^2, 0, k^2)\; 2(kP)\Big(\frac{(kP)}{k^2} - 1\Big)\Big[(\eta\epsilon) - \frac{(P\epsilon)(q\eta)}{(qP)}\Big]. \tag{11}$$

To get rid of the apparent singularity at $k^2 = 0$ we expand $A_1(\mathrm{m}_\phi^2, 0, k^2)$ for small k^2. To this purpose we use the relation between A_1 and A_2 which follows from eq.(10):

$$A_2(\mathrm{m}_\phi^2, q^2, k^2) = -\Big(1 - \frac{(kP)}{k^2}\Big)^2 A_1(\mathrm{m}_\phi^2, q^2, k^2). \tag{12}$$

Since $A_2(\mathrm{m}_\phi^2, 0, 0)$ should be finite, the amplitude $A_1(\mathrm{m}_\phi^2, q^2, k^2)$ must satisfy at $q^2 = 0$ and $k^2 = 0$ the conditions $A_1(\mathrm{m}_\phi^2, 0, 0) = \frac{\partial A_1}{\partial k^2}(\mathrm{m}_\phi^2, 0, 0) = 0$, which imply that it has to behave like k^4. Correspondingly:

$$A_2(\mathrm{m}_\phi^2, 0, 0) = -\frac{1}{2}(kP)^2\, \frac{\partial^2 A_1}{\partial^2 k^2}|_{q^2=0, k^2=0}. \tag{13}$$

In the vector meson pole approximation to eq.(7), and with ideal SU(3) mixing, only the ϕ meson intermediate state should be included. The corresponding dispersion relation in q^2 for the amplitude A_2 (which we take here as unsubtracted) involves the residue of eq.(9) at the ϕ mass shell $q^2 = p^2 = \mathrm{m}_\phi^2$, where $(kP) = 0$. At this point eq.(12) can be used to relate A_2 back to A_1, for which the on-shell residue is normalized as:

$$\overline{A}_1(\mathrm{m}_\phi^2, \mathrm{m}_\phi^2, 0) = \frac{1}{2}. \tag{14}$$

Eq.(14) easily follows from eqs.(7) and(9), using

$$< \phi(q)| \int d\vec{x}\ \theta_{00}(\vec{x}, 0)|\phi(p) >= E_p N_p, \tag{15}$$

where $E_p = \sqrt{(\vec{p})^2 + \mathrm{m}_\phi^2}$ and $N_p = (2\pi)^3\ 2E_p\ \delta(\vec{p} - \vec{q})$ is the normalization of single particle states.

Finally, by comparing eq.(11) with eq.(1), and utilizing eqs.(13) and (14) to expand $A_1(\mathrm{m}_\phi^2, 0, k^2)$, the contribution to the decay amplitude from the θ term in eq.(4) is given for small k^2 by:

$$\mathrm{G}(k^2)|_\theta = -e(\sqrt{2}\mathrm{G_F})^{1/2}\ \frac{\beta - \beta_L}{\beta_L}\ \frac{f_\phi}{\mathrm{m}_\phi^2}\ k^2, \tag{16}$$

where

$$< 0|j_\nu^{em}|\phi(q, \epsilon) >= f_\phi\ \epsilon_\nu, \tag{17}$$

with $f_\phi = 0.08\ \mathrm{GeV}^2$ from the measured $\phi \to e^+e^-$ width.

The contribution of the light quark operator of eq.(4) can be evaluated quite similarly, i.e. by starting from:

$$\epsilon_\nu(q) T_\nu(p, q, k) = \epsilon_\nu(q) \int d^4x \exp{(iqx)} < 0|\mathrm{T}\Big(j_\nu^{em}(x), \mathcal{H}_{\mathrm{SB}}(0)\Big)|\phi(p, \eta) >, \tag{18}$$

where

$$\mathcal{H}_{\mathrm{SB}} = \sum_L \mathrm{m_i \overline{q}_i q_i} \tag{19}$$

is the familiar hamiltonian density which breaks chiral $SU(3) \times SU(3)$ flavour symmetry. We expand eq.(18) compatibly with the electromagnetic current conservation:

$$T_\nu = B(p^2, q^2, k^2) \left[\eta_\nu - \frac{(q\eta)}{(qP)} P_\nu\right]. \tag{20}$$

By writing a dispersion relation in q^2, and saturating it by the ϕ, one can easily see that the matrix element $< \phi(q)|\mathcal{H}_{\mathrm{SB}}|\phi(p) >$ is the one which enters in the vector dominance appoximation to eq.(18). Specifically, the required on-shell $q^2 = p^2 = \mathrm{m}_\phi^2$ residue is normalized at $k^2 = 0$ as follows:

$$(\epsilon\eta)\overline{B}(\mathrm{m}_\phi^2, \mathrm{m}_\phi^2, 0) =< \phi(q, \epsilon)|\mathcal{H}_{\mathrm{SB}}(0)|\phi(p, \eta) > . \tag{21}$$

To make an estimate of the matrix element in eq.(21) we first write $\mathcal{H}_{\rm SB}$ in terms of SU(3) scalar densities:

$$\mathcal{H}_{\rm SB} = \sqrt{6}\mathrm{m}_0 \mathrm{U}_0 + \frac{2}{\sqrt{3}}(\overline{\mathrm{m}} - \mathrm{m}_s)\mathrm{U}_8 , \tag{22}$$

where $\mathrm{m}_0 = \frac{1}{3}(\mathrm{m_u} + \mathrm{m_d} + \mathrm{m_s})$, $\overline{\mathrm{m}} = \frac{1}{2}(\mathrm{m_u} + \mathrm{m_d})$ and (i=0,..8):

$$\mathrm{U}_i = \mathrm{Tr}\; \overline{\psi}\frac{\lambda_i}{2}\psi \tag{23}$$

with ψ the basic u, d, s quark triplet. Then, with $|\phi >= |s\bar{s} >$, the OZI rule suggests

$$< \phi|\mathrm{U}_0|\phi >= -\frac{1}{\sqrt{2}} < \phi|\mathrm{U}_8|\phi > . \tag{24}$$

Thus, taking ϕ as $\frac{1}{\sqrt{3}}\mathrm{V}_0 - \sqrt{\frac{2}{3}}\mathrm{V}_8$, and

$$< \mathrm{V}_i|\mathrm{U}_8|\mathrm{V}_j >= \mathrm{d}_{8ij}\, \mathcal{U} \tag{25}$$

with $i, j = 0, \ldots, 8$ as allowed by the nonet symmetry implied by eq.(24), and relating $\mathcal{U}$ to the mass splitting among the vector meson nonet due to U_8, one can easily derive the relation:

$$< \phi|\mathcal{H}_{\rm SB}|\phi >= \left(\mathrm{m}_\phi^2 - \mathrm{m}_{\mathrm{K}^*}^2\right) \frac{2\mathrm{m}_s}{\mathrm{m}_s - \overline{\mathrm{m}}}. \tag{26}$$

Collecting eqs.(26) and (21), and inserting them into eq.(20) we obtain after comparison with eq.(1) the quark operator contribution to the decay amplitude at $k^2 = \mathrm{m}_\mathrm{H}^2 = 0$:

$$\mathrm{G}(k^2 = 0)|_{(\bar{\mathrm{q}}\mathrm{q})} = e(\sqrt{2}\mathrm{G_F})^{1/2}\left(1 + \frac{\beta - \beta_\mathrm{L}}{\beta_\mathrm{L}}\right) \frac{\mathrm{f}_\phi}{\mathrm{m}_\phi^2} \left(\mathrm{m}_\phi^2 - \mathrm{m}_{\mathrm{K}^*}^2\right) \frac{2\mathrm{m}_s}{\mathrm{m}_s - \overline{\mathrm{m}}}. \tag{27}$$

Finally, by assembling eqs.(16) and (27) we find for $\phi \to \gamma\mathrm{H}$:

$$\mathrm{G} \cong e(\sqrt{2}\mathrm{G_F})^{1/2} \frac{f_\phi}{\mathrm{m}_\phi^2}\left[\frac{2}{9}\,\mathrm{m}_\mathrm{H}^2 + \frac{7}{9}\left(\mathrm{m}_\phi^2 - \mathrm{m}_{\mathrm{K}^*}^2\right) \frac{2\mathrm{m}_s}{\mathrm{m}_s - \overline{\mathrm{m}}}\right], \tag{28}$$

where we used $n_f = 6$ and $n_L = 3$ in eqs.(5) and (6).

Replacing eq.(28) into eq.(2) we obtain, for the Higgs boson mass $\mathrm{m_H}$ ranging from 0 to 500 MeV:

$$\mathrm{BR}(\phi \to \gamma\mathrm{H}) \cong 4.2 \cdot 10^{-9}. \tag{29}$$

Eqs.(28) and (29) depend rather weakly on the values of the current quark masses, as long as $\overline{\mathrm{m}} \ll \mathrm{m}_s$. Also, they are subject to the typical uncertainty of the vector meson dominance approach, related to the extrapolation from the vector meson mass down to the photon mass shell. On phenomenological grounds we can expect such an uncertainty to be of the order of 20-30% .

In principle we might multiply $|G|^2$ as given in eq.(28) by a Breit-Wigner factor $F_{BW}(k^2)$ normalized to unity at $k^2 = 0$, to somehow account for extrapolation effects from $k^2 = m_H{}^2 \simeq 0$ to finite values of m_H, similarly to the procedure followed in [7]. The simplest possibility for such a factor should be:

$$F_{BW}(k^2) = \frac{m_R^4\left(1 + \frac{\Gamma^2}{m_R^2}\right)}{\left(m_R^2 - k^2\right)^2 + m_R^2\Gamma^2}, \tag{30}$$

with m_R and Γ the mass and the width of a hypothetical singlet scalar resonance. In this way some enhancement of the prediction for the branching ratio given by eq.(28) can arise, depending on m_R and Γ. The size of such an effect is anyway somewhat limited by the low values of m_H considered in eq.(29).

The value of the branching ratio predicted in eq.(29) indicates that searching for a very light Higgs boson in the radiative transition $\phi \to \gamma H$ decay would require an experimental sensitivity to ϕ decays of the order of 10^{-9}. Such a small branching ratio should nevertheless be significant to the planned ϕ-factories, where sensitivities of that order of magnitude should be achieved.

The Higgs boson would manifest itself by a monochromatic photon, a signature independent of the Higgs decay modes and therefore with minimal background (which is the great advantage of two-body decays). This monochromatic photon would be accompanied for $m_H < 200$ MeV by either a photon pair from $H \to \gamma\gamma$ or by an electron-positron pair from $H \to e^+e^-$. Estimates of the corresponding widths given in [8] and [9] respectively, indicate that the latter mode should be favoured. For $m_H > 200$ MeV the dominant decays should be $H \to \mu^+\mu^-$ and $H \to \pi\pi$, as it has been recently reviewed in [7,10].

One can naively extend these considerations to the case of the J/ψ, and give a qualitative prediction for $J/\psi \to \gamma H$. To this purpose we adjoin the c quark to the sum $\sum_L m_i\bar{q}_i q_i$ in eq.(4), so that $n_L = 4$ in eqs.(5) and (6), and of course change ϕ into ψ in the various places. The estimate of the contribution from the θ term of eq.(4) remains unchanged from the case of the ϕ. The only modification comes from the quark mass term of eq.(4), which is now dominated by the large charm mass, so that we may identify in practice $\mathcal{H}_{SB} = \sum_L m_i\bar{q}_i q_i \simeq m_c\bar{c}c$. For the latter operator matrix element we may rather conservatively assume the value:

$$< \psi|m_c\bar{c}c|\psi > \simeq m_c{}^2. \tag{31}$$

Thus finally eq.(28) turns into:

$$G \simeq e(\sqrt{2}G_F)^{1/2}\frac{f_\psi}{m_\psi^2}\left[\frac{4}{25}\,k^2 + \frac{21}{25}\,m_c^2\right], \tag{32}$$

where$f_\psi = 0.8$ GeV^2 from the measured $\psi \to e^+e^-$ width. Using $m_c = 1.5$ GeV, eq.(32) would give:

$$BR(J/\psi \to \gamma H) \approx 2.5 \cdot 10^{-6}, \tag{33}$$

for the Higgs boson mass ranging between 0 and 1.5 GeV.

This estimate should be taken as simply a qualitative order of magnitude indication, since the application of the vector meson dominance approach to the J/ψ should be much more uncertain than for the previous case of the ϕ. Such an uncertainty might amount to a factor of two in the present case. Nevertheless the values in eq.(33) should still be significant, specially considering that the scale m_c used in eq.(31) could represent an underestimate. Indeed if we used there m_ψ instead of m_c, then the value in eq.(33) would be enhanced by roughly a factor of five. Also, the familiar $q\bar{q} \to \gamma H$ mechanism[11] is not applicable in the present case of the J/ψ because the one-loop QCD corrections [12,13] are overwhelming the leading order term, so that no definite statement could be done. Finally, some enhancement could arise in this case, for some values of m_H, from the Breit-Wigner factor of eq.(30).

In conclusion, we find branching ratios of the order of 10^{-9} for $\phi \to \gamma H$, and of the order of 10^{-6} for $J/\psi \to \gamma H$. Although small, these rates should be in the reach of the planned ϕ- and J/ψ-factories.

Acknowledgements

One of us (R) would like to thank the Istituto Nazionale di Fisica Nucleare, Italy, for support during this research and to acknowledge the support of the King Fahd University of Petroleum and Minerals, Dhahran, Saudi Arabia. NP gratefully acknowledges a stimulating conversation with Prof. G. Barbiellini.

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