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Decay of the Higgs boson into heavy-quarkonium states

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The exclusive production of a pair of heavy-quarkonium states V $(=\psi/J, \Upsilon, \zeta, \ldots)$ from the decay of the Higgs boson in the standard model is studied in the approximation of zero binding energy. The branching fractions could only be as large as 10^{-5} around the kinematic thresholds. The alternative process $H \to V\gamma$ is also studied and found to have branching fractions usually less than 10^{-3}.

In the standard model,[1] a neutral Higgs boson remains in the spectrum of physical particles after the spontaneous breakdown of the gauge symmetry. The unknown mass m_H of the Higgs boson can fall into a range from 300 MeV to 1 TeV without endangering the theory or contradicting present experimental limits.[2] Searching for the Higgs boson is important in understanding the symmetry-breaking mechanism. The search for the Higgs boson would be a goal of high priority in future high-energy accelerators. However, the experimental difficulty lies in the absence of techniques for tagging the heavy particles which are the dominant decay products.[3,4]

In this paper we study the Higgs-boson decays into heavy vector quarkonium states V $(=\psi/J$, Υ, and the presumed t-quarkonium $\zeta)$, namely, the processes $H \to VV$ and also $H \to V\gamma$. These produced quarkonia could easily be identified via their subsequent decays into lepton pairs. The rates are usually suppressed by the binding effects.

Firstly, we calculate the partial width for $\Gamma(H \to VV)$. Figure 1 illustrates the lowest-order diagrams for the process $H \to VV$. Within the zero-binding-energy approximation,[5] $m_Q = \frac{1}{2} m_V$, the momenta of the heavy quark Q and the heavy antiquark $\bar{Q}$ are equal. The amplitude T for producing a pair of V of the momenta p and p' and polarizations e and e' is given as

$$T = \tfrac{4}{3}(128\alpha_S \pi \phi_0^2/m_H^4)(\sqrt{2}G_F m_Q^2)^{1/2} \times [2p\cdot e' e\cdot p' - (m_H^2 - 2m_V^2) e\cdot e'] \ . \quad (1)$$

Here the factor $\frac{4}{3}$ comes from the color factor. The wave functions at the origin are measured through the leptonic decay width Γ_{ee}:

$$\phi_0^2 = m_V^2 \Gamma_{ee}/16\pi\alpha^2 e_Q^2 \ . \quad (2)$$

The values of ϕ_0^2 for ψ/J and Υ are, therefore,

$$\phi_0^2(\psi/J) = 0.04\ \text{GeV}^3 \ , \qquad \phi_0^2(\Upsilon) = 0.4\ \text{GeV}^3 \ , \quad (3)$$

respectively. We also set the strong coupling constant $\alpha_S = 0.3$ in our calculations. For the $\zeta(t\bar{t})$ production, we assume $m_\zeta = 40$ GeV and approximate the wave-function value ϕ_0^2 from the calculation based on a Coulomb-type binding potential

$$\phi_0^2 = \alpha_S^3 m_V^3/27\pi \ . \quad (4)$$

The partial decay width $\Gamma(H \to VV)$ is

$$\Gamma(H \to VV) = \frac{1}{2}\frac{1}{2m_H}\frac{1}{8\pi}\sum_{\text{spins}} |T|^2 \left(1 - \frac{4m_V^2}{m_H^2}\right)^{1/2} \quad (5)$$

$$= \frac{2^{14}\pi\alpha_S^2\sqrt{2}G_F m_Q^2 m_H}{9}\left(\frac{\phi_0^2}{m_H^3}\right)^2 \left(1 - \frac{4m_V^2}{m_H^2} + \frac{6m_V^4}{m_H^4}\right)\left(1 - \frac{4m_V^2}{m_H^2}\right)^{1/2} \ . \quad (6)$$

The nature of identical particles in the final state gives rise to the first factor $\frac{1}{2}$ in Eq. (5). In comparison to the dominant decay channel $H \to Q\bar{Q}$,

$$\Gamma(H \to Q\bar{Q}) = 3(\sqrt{2}G_F m_Q^2 m_H/8\pi)(1 - m_V^2/m_H^2)^{3/2} \ , \quad (7)$$

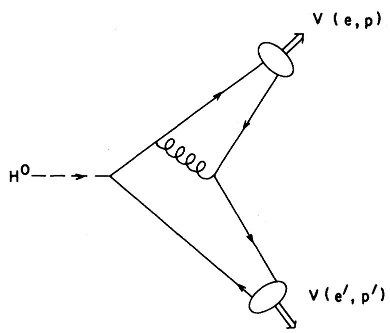

FIG. 1. Lowest-order QCD diagrams for $H \to VV$. Three other similar diagrams are not shown.

we obtain the branching ratio

$$\frac{\Gamma(H\to VV)}{\Gamma(H\to Q\bar{Q})}=\frac{2^{17}\pi^2\alpha_S{}^2\phi_0{}^4}{27m_H{}^6}\left[1-\frac{4m_V{}^2}{m_H{}^2}+\frac{6m_V{}^4}{m_H{}^4}\right]\times\frac{(1-4m_V{}^2/m_H{}^2)^{1/2}}{(1-m_V{}^2/m_H{}^2)^{3/2}} \quad . \qquad (8)$$

Our results are shown in Fig. 2. The branching fractions could only be as large as 10^{-5} around the kinematic thresholds. They might be too rare for detection even though their identifications are simpler than other dominant channels.

Now we study another process, $H \to V\gamma$, the inverse Wilczek[6] mechanism. The rate is bigger than the previous one as only one binding vertex $VQ\bar{Q}$ is involved. The amplitude is

$$T=\sqrt{3}4e_Q e(G_F m_Q \phi_0{}^2/\sqrt{2})^{1/2}(m_H{}^2-m_V{}^2)^{-1}\times[2p_\gamma\cdot e_V p_V\cdot e_\gamma-(m_H{}^2-m_V{}^2)e_\gamma\cdot e_V] \quad . \qquad (9)$$

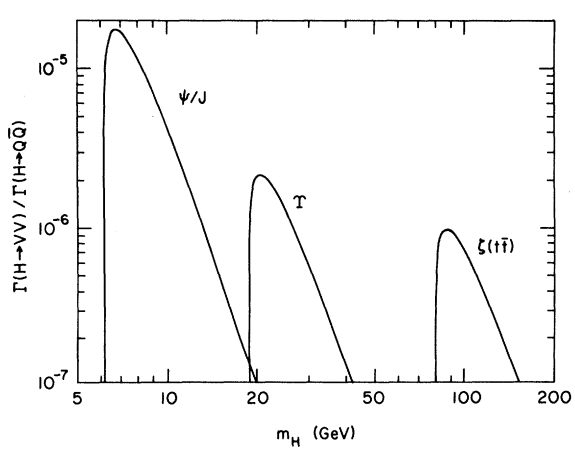

FIG. 2. Ratio of $\Gamma(H\to VV)$ to $\Gamma(H\to Q\bar{Q})$ vs m_H for the cases $V=\psi/J$, Υ, or ζ, assuming $m_\zeta = 40$ GeV.

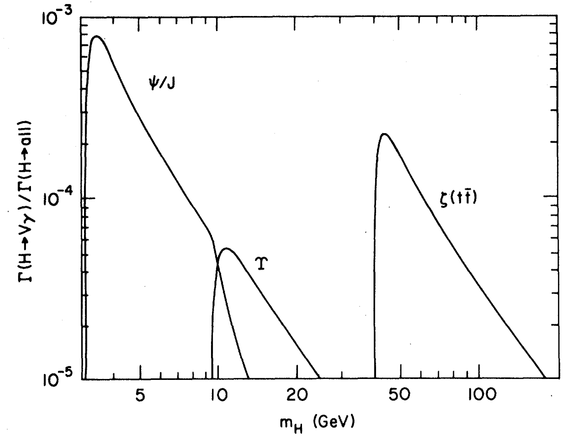

FIG. 3. Branching fraction of $H\to V\gamma$ vs m_H for the cases $V=\psi/J$, Υ, and ζ, assuming $m_\zeta=40$ GeV.

Here e_γ and e_V are the polarizations and p_γ and p_V are the momenta of γ and V, respectively. The partial width is

$$\Gamma(H\to V\gamma)=6\sqrt{2}G_F m_V \alpha e_Q{}^2(m_H{}^2-m_V{}^2)(\phi_0{}^2/m_H{}^3) \quad . \qquad (10)$$

Figure 3 illustrates the branching fractions versus m_H for the cases $V=\psi/J$, Υ, or the hypothetical t-quarkonium of mass $m_\zeta=40$ GeV. The branching fractions lie in a range below 10^{-3}. It should be noticed that the zero-binding approximation fails for the soft-photon emission when $m_H \approx m_V$, and the results are overestimated.

Our conclusion is that the quarkonia are produced quite rarely from the decay of the Higgs boson. They can only be observed in a copious source of the Higgs bosons.

Note added. The decay of the Higgs boson into quarkonium states has been studied previously by M. Bander and A. Soni [Phys. Lett. 82B, 411 (1979)]. Their approach to the process $H\to VV$ is different from ours. They do not include the gluon line in Fig. 1 so that the quark lines cannot all be maintained on shell. This requires knowledge of the off-shell behavior of the form factor, which is not well understood at the present time. By taking a rather mild form factor, they obtain a much larger branching fraction $B(H\to VV)$ than ours, which we believe, is based on better known parameters.

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