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HIGGS PARTICLE PRODUCTION BY $Z \to H\gamma$ †

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ABSTRACT

The rate for the decay of a Z-boson into a Higgs boson and monochromatic photon is computed to leading order in the standard $SU(2) \times U(1)$ gauge theory. The coupling has contributions from fermion and W-boson loops. The W-boson loop dominates unless the number of heavy fermion generations exceeds six. The branching ratio computed from the W-boson loop contribution, $B(Z \to H\gamma)$, is approximately $2 \times 10^{-6} \left(1 - \frac{M_H^2}{M_Z^2}\right)^3$.

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In its simplest form the $SU(2) \times U(1)$ gauge theory of weak and electromagnetic interactions requires the existence of a single, massive, neutral Higgs boson.[1] This scalar particle is the remnant of the mechanism responsible for the spontaneous symmetry breaking which gives masses to the physical particles other than the photon. The mass of the Higgs particle itself is not determined by the theory, although there are arguments which suggest it should be more than a few GeV.[2] The discovery of the Higgs boson would be a remarkable verification of the original proposal of Weinberg and Salam.[3] Finding this particle will be a formidable task, both because its production is suppressed by its small couplings to light fermions $(M_F << M_W)$ and because its decays will be into multiparticle states.

It is natural to consider Z decays as a source for Higgs bosons, since it is only at the Z resonance that future very high energy e^+e^- machines will have a significant event rate.[4] At the peak of the resonance, the cross section for Z production is equal to the usual point electromagnetic cross section $(4\pi\alpha^2/3s)$, multiplied by the factor $BR\,(Z \to e^+e^-) \times 9/\alpha^2$. A reasonable estimate* of the $Z \to e^+e^-$ branching ratio is 3%. Thus R, the ratio of the cross section to the usual point cross section would be about 5100. For a luminosity of $10^{32}\ cm^{-2}\ sec^{-1}$ this gives an event rate of 4×10^5/day. Besides the rate of production, which depends on the luminosity of a machine not yet designed, the signature of the proposed process is also extremely important.

* Neglecting phase space effects, the branching ratio $B(Z \to \mu^+\mu^-) \cong 1/11N$, where N is the number of fermion generations with thresholds below M_Z.

We have calculated in the conventional $SU(2) \times U(1)$ theory the partial decay rate for $Z \to H + \gamma$. This coupling vanishes in lowest order since at that level the photon couples only to charged particles. Thus the process occurs only through higher order diagrams. Fermion loops and W-boson loops both contribute. The fermion loops are suppressed by a factor M_F^2/M_W^2 when $M_F \ll M_W$. For $M_F \gtrsim M_W$ this suppression disappears but, even in this case, the W-boson loop contribution is much larger than the contribution from a single fermion species. The fermion loop contribution is only important if there are many quark and/or lepton flavors at least as heavy as the W-boson. The rate for $Z \to H\gamma$ would then "count" the number of heavy flavors.†

The rate we have calculated for $Z \to H\gamma$ from the W-boson loop contribution alone corresponds to a branching ratio of about 2×10^{-6} for $M_H \ll M_Z$, making its observation difficult but competitive with any other proposed means. The most attractive decay previously suggested is $Z \to H\,\mu^+\mu^-$ where the muon pair arises from the decay of a virtual Z.[6,7] This process has a rather good signature: the missing mass recoiling against the muon pair. The branching ratio is quite large, 7×10^{-5}, for a 10 GeV Higgs boson, but drops rapidly with increasing Higgs mass, so that for a Higgs boson of 40 GeV, the branching ratio is about 4×10^{-6}. We expect that the cleaner signature of $Z \to H\gamma$ may compensate for its lower branching ratio.

† In a similar way the hadronic decays of the Higgs via two gluons count the number of quark flavors heavier than the Higgs - c.f. Ref. 5.

The amplitude for the decay of a Z with momentum p and polarization vector η into a Higgs boson and a photon with momentum k and polarization ε is

$$\mathcal{M} = \varepsilon^\mu \mathcal{M}_{\mu\nu} \eta^\nu \tag{1}$$

where electromagnetic gauge invariance requires $\mathcal{M}_{\mu\nu}$ to be of the form

$$\mathcal{M}_{\mu\nu} = (k_\nu p_\mu - k \cdot p\, g_{\mu\nu})\mathcal{A}. \tag{2}$$

The decay width is then

$$\Gamma(Z \to H\gamma) = \frac{E_\gamma^3 \mathcal{A}^2}{12\pi} , \tag{3}$$

where E_γ is the photon energy in the Z rest frame.

For the standard form of $SU(2) \times U(1)$ considered here the Z and W masses are related by

$$M_Z^{\,2} = M_W^{\,2}/\cos^2\theta_w , \tag{4}$$

and the gauge coupling constant, g, and the electric charge are related by

$$g^2 = e^2/\sin^2\theta_w , \tag{5}$$

where θ_w is the Weinberg angle. Experimentally, $\sin^2\theta_w \cong 0.20$. A convenient reference for Z decays is the partial width into muons:

$$\Gamma(Z \to \mu^+\mu^-) = \frac{(g_V^{\,2} + g_A^{\,2})M_Z}{12\pi} . \tag{6}$$

In the standard model,

$$g_V = g(-1 + 4\sin^2\theta_w)/(4\cos\theta_w),$$

$$g_A = -g/(4\cos\theta_w)$$

which yields a width into muons of about 94 MeV.

We begin our sketch of the calculation of the $Z \to H\gamma$ decay rate with a discussion of the fermion loop contribution. Since the Higgs particle is C-even and the photon is C-odd, only the C-odd (vector) coupling of the Z contributes. Because of the experimental value of $\sin^2\theta_w$, the vector couplings are much smaller than the axial ones. This significantly reduces the importance of the fermion loop diagrams. Furthermore, since the Higgs boson coupling to fermions is proportional to their mass, fermions like e, μ, u, d, s, τ, c, b, which are much lighter than the Z, contribute only slightly.

The contribution of each fermion may be written conveniently by extracting an overall factor from $\mathcal{a}$

$$\mathcal{a} = \frac{eg^2\pi^2}{(2\pi)^4 M_W} A, \tag{7}$$

so that A is dimensionless and the decay rate relative to muon pairs is

$$\frac{\Gamma(Z \to H\gamma)}{\Gamma(Z \to \mu^+\mu^-)} = \frac{\alpha^2}{\pi^2\sin^2\theta_w}\left(\frac{E_\gamma}{M_Z}\right)^3 \frac{A^2}{[1 + (1 - 4\sin^2\theta_w)^2]}$$

$$\simeq 2.60 \times 10^{-5}\left(\frac{E_\gamma}{M_Z}\right)^3 A^2. \tag{8}$$

The loop integration is handled easily using dimensional regularization yielding a result of the gauge invariant form given in Eq. (2). If the fermion has charge Q_F (in units of the proton charge) and a third component of weak isospin T_F^{3L}, the contribution is

$$A_F = \frac{-2\,Q_F\,(T_F^{3L} - 2\,Q_F\sin^2\theta_w)}{\cos\theta_w} \times \int_0^1 dx \int_0^{1-x} dy \; \frac{(4xy - 1)\,M_F^2}{M_F^2 - y(1-y)\,M_Z^2 + xy(M_Z^2 - M_H^2)}. \tag{9}$$

In the light fermion limit, A_F is approximately of order M_F^2/M_W^2 and can be ignored. For very large fermion masses, A_F approaches the value

$$A_F \underset{M_F\to\infty}{\sim} \frac{2}{3}\,\frac{Q_F(T_F^{3L} - 2\,Q_F\sin^2\theta_w)}{\cos\theta_w}. \tag{10}$$

For the possibilities of a heavy charged lepton ($Q_F = -1$, $T_F^{3L} = -\frac{1}{2}$), a heavy up quark ($Q_F = 2/3$, $T_F^{3L} = \frac{1}{2}$), and a heavy down quark ($Q_F = -1/3$, $T_F^{3L} = -\frac{1}{2}$), this yields values of A_F of about 0.07, 0.12, and 0.09 respectively. Of course all contributions must be added coherently. Suppose, for example, there is a complete generation of heavy fermions with masses large compared to the Z mass. This would give a contribution

$$A_{\text{Fermion Generation}} = \frac{4}{3\cos\theta_w}\left(1 - \frac{8}{3}\sin^2\theta_w\right) \tag{11}$$

$$\cong 0.7\,.$$

Using Eq. (8) we see that a single generation of heavy fermions would by itself give a branching ratio of at most 5×10^{-8} which is too small to be useful.

The W-boson loop contribution is precisely defined within the $SU(2) \times U(1)$ model. It must be finite since, as seen from Eq. (2), any counterterm would be constructed from the product of the Higgs field and the contraction of the Z and photon field strengths, a combination which would destroy renormalizability. We have done the calculation in both the unitary and 't Hooft-Feynman gauges,[8] using dimensional regularization. In either gauge there are many internal consistency checks. The finiteness of the result and the gauge invariance of the finite answer appear only after all diagrams are added. Both calculations are quite tedious. In the 't Hooft-Feynman gauge, the vector propagator has only $g_{\mu\nu}$ in the numerator so each diagram is relatively simple. However, in this gauge, there are contributions from ghost Higgs particles and from Fadeev-Popov ghost loops. As a consequence, there are about thirty graphs to evaluate. In the unitary gauge there are only three diagrams to consider, but this reduction is compensated by the enormously greater complexity arising from the presence of two terms in each vector propagator. The calculation in a general ξ gauge[9] would combine the difficulties of both the unitary and the 't Hooft-Feynman gauges and has not been attempted.

The result of these calculations is

$$A_W = -4 \cos\theta_W M_W^2 \times \left\{ \int_0^1 dx \int_0^{1-x} dy \, \frac{(3-\tan^2\theta_W) + xy\left[\left(1+\frac{M_H^2}{2M_W^2}\right)\tan^2\theta_W - \left(5+\frac{M_H^2}{2M_W^2}\right)\right]}{M_W^2 - y(1-y)M_Z^2 + xy\,(M_Z^2 - M_H^2)} \right\} . \tag{12}$$

The integral is non-singular and a very good approximation may be obtained by expanding the denominator,

$$\int_0^1 dx \int_0^{1-x} dy \, \frac{M_W^2}{M_W^2 - y(1-y)\,M_Z^2 + xy(M_Z^2 - M_H^2)} = \frac{1}{2} + \frac{1}{12}\frac{M_Z^2}{M_W^2} - \frac{1}{24}\left(\frac{M_Z^2 - M_H^2}{M_W^2}\right) + \ldots , \tag{13a}$$

$$\int_0^1 dx \int_0^{1-x} dy \, \frac{M_W^2 yx}{M_W^2 - y(1-y)M_Z^2 + xy(M_Z^2 - M_H^2)} = \frac{1}{24} + \frac{1}{120}\frac{M_Z^2}{M_W^2} - \frac{1}{180}\frac{(M_Z^2 - M_H^2)}{M_W^2} + \ldots . \tag{13b}$$

In this way we obtain an adequate representation for A_W:

$$A_W = -\left[4.9 + 0.3\,(M_H^2/M_W^2)\right] . \tag{14}$$

Inserting this into Eq. (8), and neglecting possible contributions from heavy fermion loops, we obtain the approximate form

$$\frac{\Gamma(Z \to H\gamma)}{\Gamma(Z \to \mu^+\mu^-)} \approx 7.8 \times 10^{-5} \left(1 - \frac{M_H^2}{M_Z^2}\right)^3 \left(1 + 0.17 \frac{M_H^2}{M_Z^2}\right). \quad (15)$$

The results are shown in Fig. 1, along with the corresponding results for the process $Z \to H \mu^+\mu^-$.

If there are fermion flavors as heavy as or heavier than the Z-boson, the fermion and W boson loop contributions to the decay amplitude add coherently. Comparison of Eq. (11) and Eq. (14) shows that the interference is destructive for $\sin^2 \theta_W < 3/8$.

Several other means have been suggested for searching for the Higgs boson. One is the process $e^+e^- \to Z + H$, which is similar to the alternative mentioned above, but requires that the incident energy be high enough to produce a Higgs particle and a real Z. An additional difficulty is finding a suitable signature for this process. The suggestion of Wilczek[4] to look for $V \to H\gamma$ suffers from the problem of rate, as do the others. Its utility is also limited by the mass range which can be scanned. The recent proposal of Gaemers and Gounaris[10] to produce Higgs particles by bremsstrahlung from heavy fermions produced in e^+e^- collision is weakest on the question of signature.

It is clear that the difficulty of finding the Higgs is not due just to the low production rates expected, but also to the lack of a clear signature in its decay. This is the special virtue of the decay $Z \to H\gamma$, and to a somewhat lesser extent of $Z \to H\mu^+\mu^-$: the signature is found in the system recoiling against the Higgs particle.

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5. F. Wilczek, Phys. Rev. Lett. 39, 1304 (1977).
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7. J. D. Bjorken, Weak Interaction Theory and Neutral Currents, Proceedings of the SLAC Summer Institute, 1976, SLAC Report #198. Note that there is a typographical error in Eq. (4.30). There should be an additional factor π in the denominator. The curve in Fig. 11 is correctly calculated, however.
8. J. Ellis et al., Ref. 1, have calculated the decay $H \to \gamma\gamma$ in the 't Hooft-Feynman gauge. The diagrams and vertices displayed in that reference are a subset of those required for our process. The appropriate Feynman rules have been extracted from K. Fujikawa, B. W. Lee, and A. I. Sanda, Phys. Rev. D6, 2923 (1972), Eqs. (C20) - (C29).

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10. K. J. F. Gaemers and G. J. Gounaris, Phys. Lett. 77B, 379 (1978).

FIGURE CAPTIONS

Fig. 1: Comparison of the decay rates for $Z \to H\gamma$ and $Z \to H\,\mu^+\mu^-$. The curves show the ratios $\Gamma(Z \to H\gamma)/\Gamma(Z \to \mu^+\mu^-)$ (solid line) and $\Gamma(Z \to H\,\mu^+\mu^-)/\Gamma(Z \to \mu^+\mu^-)$ (dashed line) as a function of the mass ratio M_H/M_Z. The curves are computed for $\sin^2\theta_W = .2$ so that $M_Z = 94$ GeV.

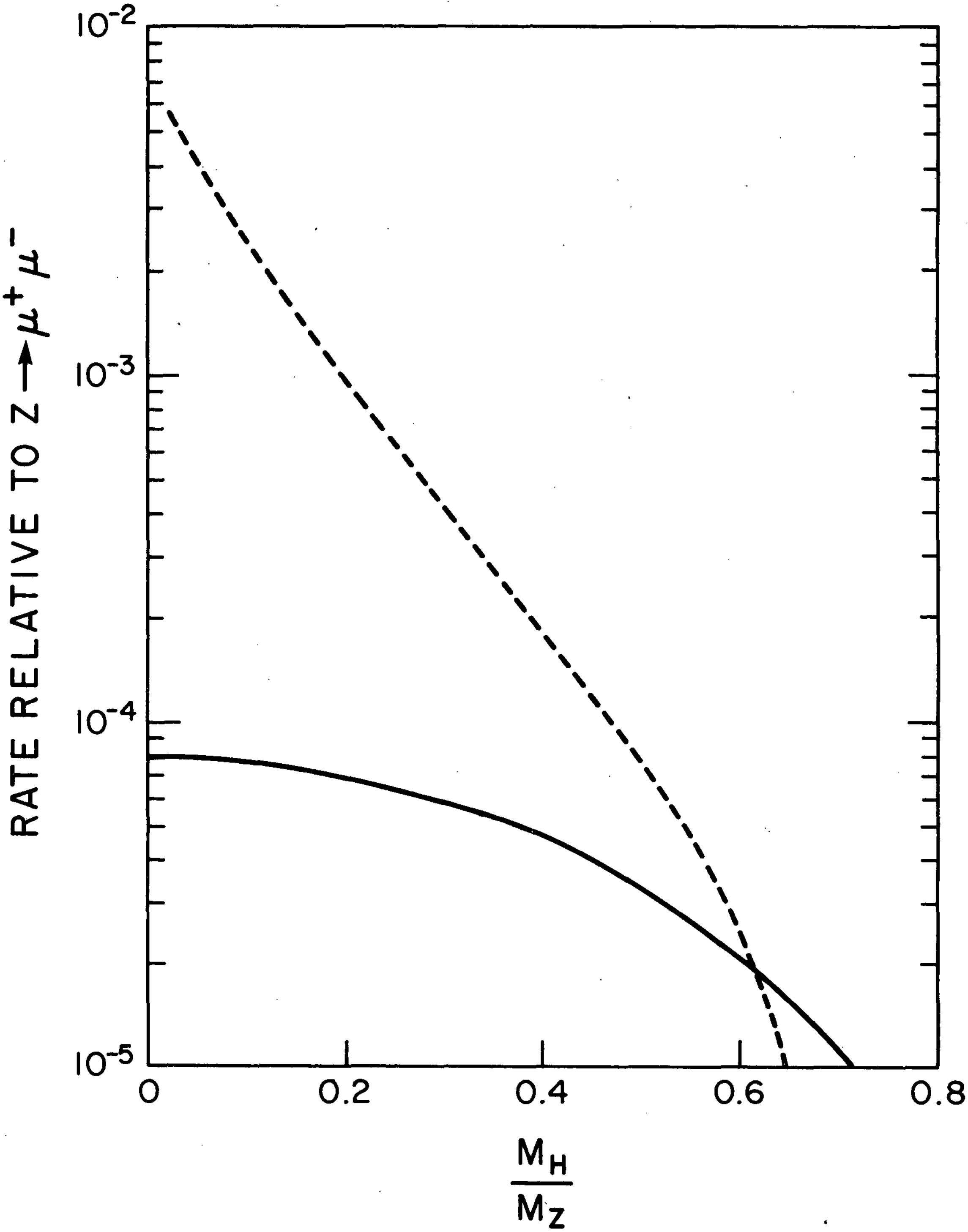

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Fig. 1

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