Decay of Z into three pseudoscalar bosons

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We consider the decay of the Z boson into three pseudoscalar bosons in a general two-Higgs-doublet model. Assuming m_A to be very small, and that of the two physical neutral scalar bosons h_1 and h_2, A only couples to Z through h_1, we find the $Z\to AAA$ branching fraction to be negligible for moderate values of $\tan\beta \equiv v_2/v_1$, if there is no $\lambda_5(\Phi_1^\dagger\Phi_2)^2+$H.c. term in the Higgs potential; otherwise there is no absolute bound but very large quartic couplings (beyond the validity of perturbation theory) are needed for it to be observable. [S0556-2821(98)04103-4]

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If the standard SU(2)×U(1) electroweak gauge model is extended to include two scalar doublets, there will be a neutral pseudoscalar boson A whose mass may be small. In that case, the decay of the Z boson into 3 A's may not be negligible. This process was first studied [1] in a specific model [2]. It was then discussed [3] in a more general context. More recently, it has been shown [4] that there is a lower bound on m_A of about 60 GeV in the minimal supersymmetric standard model (MSSM), hence the decay $Z\to AAA$ is only of interest for models with two scalar doublets of a more general structure. Even in the context of supersymmetry, this is possible [5] if there exists an additional U(1) gauge factor at the TeV scale.

In this paper we consider a general two-Higgs-doublet model and identify the conditions for which the decay $Z\to AAA$ may be enhanced, despite the nonobservation of $e^+e^-\to h+A$, where h is either one of the two neutral scalar bosons of the model. We will show that in principle this decay is limited only by the scalar coupling $\lambda_1-\lambda_2$ as defined below. However, if $\lambda_5=0$, which is true in a large class of models [6], then it may be bounded as discussed below.

Let the Higgs potential V for two SU(2)×U(1) scalar doublets $\Phi_{1,2}=(\phi_{1,2}^+,\phi_{1,2}^0)$ be given by

$$V=m_1^2\Phi_1^\dagger\Phi_1+m_2^2\Phi_2^\dagger\Phi_2+m_{12}^2(\Phi_1^\dagger\Phi_2+\Phi_2^\dagger\Phi_1)+\frac{1}{2}\lambda_1(\Phi_1^\dagger\Phi_1)^2+\frac{1}{2}\lambda_2(\Phi_2^\dagger\Phi_2)^2+\lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2)+\lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1)+\frac{1}{2}\lambda_5(\Phi_1^\dagger\Phi_2)^2+\frac{1}{2}\lambda_5^*(\Phi_2^\dagger\Phi_1)^2, \tag{1}$$

where the discrete symmetry $\Phi_1\to\Phi_1$ and $\Phi_2\to-\Phi_2$ is only broken softly by the m_{12}^2 term. Assume λ_5 to be real for simplicity. Define $\tan\beta\equiv v_2/v_1$ as is customary, where $v_{1,2}=\langle\phi_{1,2}^0\rangle$ are the usual two nonzero vacuum expectation values. The pseudoscalar neutral Higgs boson is then

$$A=\sqrt{2}(\sin\beta\ \mathrm{Im}\ \phi_1^0-\cos\beta\ \mathrm{Im}\ \phi_2^0), \tag{2}$$

with mass given by

$$m_A^2=-m_{12}^2(\tan\beta+\cot\beta)-2\lambda_5 v^2, \tag{3}$$

where $v^2\equiv v_1^2+v_2^2$, and the charged Higgs boson is

$$h^\pm=\sin\beta\phi_1^\pm-\cos\beta\phi_2^\pm, \tag{4}$$

with

$$m_{h^\pm}^2=m_A^2+(\lambda_5-\lambda_4)v^2. \tag{5}$$

To get the maximum $Z\to AAA$ rate, we let $m_A=0$, i.e.

$$m_{12}^2=-2\lambda_5 v^2\sin\beta\cos\beta. \tag{6}$$

Then the mass-squared matrix spanning the two neutral scalar Higgs bosons $\sqrt{2}\mathrm{Re}\phi_{1,2}^0$ is given by

$$\mathcal{M}^2=2v^2\begin{bmatrix}\lambda_1\cos^2\beta+\lambda_5\sin^2\beta & (\lambda_3+\lambda_4)\sin\beta\cos\beta\\ (\lambda_3+\lambda_4)\sin\beta\cos\beta & \lambda_2\sin^2\beta+\lambda_5\cos^2\beta\end{bmatrix}. \tag{7}$$

Consider now the following two linear combinations:

$$h_1=\sqrt{2}(\sin\beta\ \mathrm{Re}\phi_1^0-\cos\beta\,\mathrm{Re}\phi_2^0), \tag{8}$$

$$h_2=\sqrt{2}(\cos\beta\,\mathrm{Re}\phi_1^0+\sin\beta\,\mathrm{Re}\phi_2^0). \tag{9}$$

It is well known that h_1 couples to AZ but not ZZ, whereas h_2 couples to ZZ but not AZ. However, the process $e^+e^-\to h+A$ is in general possible because h will normally have a h_1 component, thereby putting a constraint on m_A if kinematically allowed. For our purpose, we will require h_1 and h_2 to be mass eigenstates, in which case m_A is uncontrained by the nonobservation of $e^+e^-\to h+A$ even if m_2 is small, as long as m_1 is larger than the e^+e^- center-of-mass energy. This allows us to have the maximum effective coupling of Z to AAA as shown below.

The requirement that h_1 and h_2 be mass eigenstates leads to the condition

$$\lambda_2\sin^2\beta-\lambda_1\cos^2\beta+(\lambda_3+\lambda_4+\lambda_5)(\cos^2\beta-\sin^2\beta)=0. \tag{10}$$

As a result, the masses of $h_{1,2}$ are given by

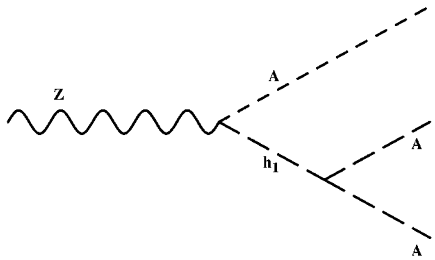

FIG. 1. One of 3 diagrams for the decay $Z \to AAA$. The other 2 are obvious permutations.

$$m_1^2 = [\lambda_1 \cos^2\beta + \lambda_2 \sin^2\beta + \lambda_5 - \lambda_3 - \lambda_4] v^2, \quad (11)$$

$$m_2^2 = [\lambda_1 \cos^2\beta + \lambda_2 \sin^2\beta + \lambda_5 + \lambda_3 + \lambda_4] v^2. \quad (12)$$

Note that in the MSSM, Eq. (10) cannot be satisfied in the presence of radiative corrections.

We now extract the h_1AA coupling from Eq. (1), using Eqs. (2) and (8). We find it to be given by

$$\frac{\sin 2\beta}{2\sqrt{2}}(\lambda_1 - \lambda_2) v, \quad (13)$$

where Eq. (10) has been used. As a function of β, this expression is obviously maximized at $\sin 2\beta = \pm 1$. On the other hand, our conditions so far do not limit the combination $\lambda_1 - \lambda_2$, hence there is no absolute bound on $Z \to AAA$ in this general case.

Let us consider the case $\lambda_5 = 0$. This is natural in a large class of models where the two Higgs doublets are remnants [6] of a gauge model larger than the standard model such that they are distinguishable under the larger symmetry. In that case, we have

$$m_1^2 = 2(\lambda_1 - \lambda_3 - \lambda_4) v^2 \cos^2\beta = 2(\lambda_2 - \lambda_3 - \lambda_4) v^2 \sin^2\beta, \quad (14)$$

and we can rewrite (13) as

$$-\frac{m_1^2}{v\sqrt{2}} \cot 2\beta. \quad (15)$$

The above expression appears to be unbounded as $\sin 2\beta \to 0$. However, that would require very large quartic scalar couplings. This can be seen two ways. First, since (15) is equal to (13), we need an extremely large value of $\lambda_1 - \lambda_2$. Second, from Eq. (14), we see also that if $\sin\beta$ is small, then $\lambda_2 - \lambda_3 - \lambda_4$ has to be big, and if $\cos\beta$ is small, then $\lambda_1 - \lambda_3 - \lambda_4$ has to be big. Thus we will choose moderate values of $\tan\beta$ in (15) for the following discussion.

In Fig. 1 we show the diagram for the decay $Z \to AAA$ with an intermediate virtual h_1. To maximize this rate, we minimize m_1 to be just above the maximum experimental e^+e^- center-of-mass energy, which is 172 GeV up to now but will soon be 183 GeV. As for h_2, it interacts exactly as the one Higgs boson of the standard model, from which we have the experimental limit [7] of $m_2 > 65$ GeV. However, m_2 is not directly involved in the h_1AA coupling here. Note also that λ_4 by itself must be large and negative so that $m_{h^\pm}$ of Eq. (5) can be greater than $m_t - m_b$ for $m_A = 0$, so as to prevent the decay $t \to b + h^+$. This condition is not satisfied in the MSSM where $\lambda_4 = -g_2^2/2$, hence $m_A = 0$ is not allowed there [4].

Assuming $\lambda_5 = 0$ and using Eq. (15) with $m_1 = 180$ GeV and $|\cot 2\beta| = 1$ (i.e. $\tan\beta = 0.4$ or 2.4), we now calculate the $Z \to AAA$ decay rate, following Ref. [1]. The amplitude is given by

$$\mathcal{M} = g_Z \frac{m_1^2\sqrt{2}}{v}\left[\frac{\epsilon\cdot k_1}{(p-k_1)^2 - m_1^2} + \frac{\epsilon\cdot k_2}{(p-k_2)^2 - m_1^2} + \frac{\epsilon\cdot k_3}{(p-k_3)^2 - m_1^2}\right], \quad (16)$$

where $g_Z = e/\sin\theta_W \cos\theta_W$, p is the four-momentum of the Z boson, and $k_{1,2,3}$ are those of the A's. The effective coupling used in Ref. [1] is now determined to be

$$\lambda_{\rm eff} = \frac{m_1^2\sqrt{2}}{v^2} \simeq 1.5. \quad (17)$$

Using the estimate of Ref. [1], this $Z \to AAA$ rate is then about 1.0×10^{-7} GeV. Hence its branching fraction is about 4×10^{-8} which is clearly negligible. To obtain a branching fraction of 10^{-6}, we need $\cot 2\beta = 5$ (i.e. $\tan\beta = 0.1$ or 10). In this case, either $\lambda_1 - \lambda_3 - \lambda_4$ or $\lambda_2 - \lambda_3 - \lambda_4$ in Eq. (14) has to be about 53.5. If $\lambda_5 \neq 0$, then we cannot use Eqs. (14) and (15), but Eq. (13) is still valid. To obtain a branching fraction of 10^{-6}, we will then need $|\lambda_1 - \lambda_2|$ to be about 53.5. Thus in both scenarios, one or more quartic scalar couplings have to be very large and beyond the validity of perturbation theory.

If h_1 and h_2 are not exact mass eigenstates, then there is an additional contribution from $h_1 - h_2$ mixing which is necessarily very small from the constraint of experimental data if m_2 is below 172 GeV. The h_2AA coupling is given by

$$\frac{v}{\sqrt{2}}\left(\frac{m_2^2}{2v^2} - 2\lambda_5[1 - \sin^2\beta \cos^2\beta]\right). \quad (18)$$

If $\lambda_5 = 0$, this expression is bounded independent of $\tan\beta$ and the overall contribution (including the small $h_1 - h_2$ mixing) is negligible. If $\lambda_5 \neq 0$, then its value has to be huge for the process to be observable.

The reason that $\Gamma(Z \to AAA)$ is so small is twofold. One is that with the higher energy reached by the CERN e^+e^- collider LEP2, the nonobservation of $Z \to h + A$ forces m_1 to be much greater than M_Z. The other is that for $m_1 \gg M_Z$, the leading term in $\mathcal{M}$ vanishes because $\epsilon \cdot (k_1 + k_2 + k_3) = 0$, resulting in a very severe suppression factor [1]. Our conclusion is that the decay $Z \to AAA$ is not likely to be observable in a general two-Higgs-doublet model with parameters in the perturbative regime.

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