Measuring the Spin of the Higgs Bosons

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Abstract

We summarize two procedures for determining spin and parity of the Higgs boson(s) in the Standard Model and related extensions unambiguously in a model independent way. One is to study the excitation curve near threshold and the angular distribution in Higgs–strahlung, $e^+e^- \to ZH$ and the other to exploit the decay, $H \to ZZ$. In the latter process with a Higgs mass above the on-shell ZZ threshold, a complete model–independent analysis can be performed only by using additional angular correlation effects in gluon–gluon fusion at the LHC and $\gamma\gamma$ fusion at linear colliders. In the intermediate mass range, in which the Higgs boson decays into pairs of real and virtual Z bosons, threshold effects and angular correlations, parallel to Higgs-strahlung, can be adopted to determine spin and parity.

1 Introduction

The Higgs boson in the Standard Model (SM) must necessarily be a scalar particle, assigned the external quantum numbers $\mathcal{J}^{\mathcal{PC}} = 0^{++}$; extended models such as $\mathcal{CP}$–invariant supersymmetric theories also contain these pure scalar states. The assignment of the quantum numbers invites investigating experimental opportunities to identify spin and parity of the Higgs state at future high–energy colliders, see Refs.[1] and [2]. The determination of the parity and the parity mixing of spinless Higgs bosons have been extensively investigated in Refs.[3, 4].

In this talk we summarize methods for identifying the spinless nature and the positive parity of the Higgs boson unambiguously in Higgs–strahlung [1] and the decay process [2]:

$$e^+e^- \to ZH \qquad \text{and} \qquad H \to ZZ \tag{1}$$

These processes involve clean $\mu^+\mu^-$ and e^+e^- subsequent decay channels of the Z bosons for isolating the signal processes from the background and allowing a complete reconstruction of the kinematical configuration with good precision. Note that while the dominant decay mode for Higgs masses below ~ 140 GeV is the $b\bar{b}$ decay channel, the ZZ mode, one of the vector bosons being virtual below the threshold for two real Z bosons, becomes leading for higher masses next to the WW decay channel.

2 General description

Without loss of generality, the Higgs state H can be assumed to be emitted from the Z–boson line. Were it emitted from the lepton line, the required Hee coupling would be so large that the state could have been detected as a resonance at LEP, $e^+e^- \to H(\gamma)$. The most general form of the HZZ vertex is given by the expression

$$\mathcal{T} = \frac{g_W M_Z}{\cos\theta_W}\, T_{\mu\nu\beta_1\ldots\beta_\mathcal{J}}\, Z_1^\mu\, Z_2^\nu\, H^{\beta_1\ldots\beta_\mathcal{J}} \tag{2}$$

While Z_1^μ and Z_2^ν are the usual spin–1 polarization vectors, the spin–$\mathcal{J}$ polarization tensor $H^{\beta_1\ldots\beta_\mathcal{J}}$, constructed from products of suitably chosen polarization vectors, is symmetric, traceless and orthogonal to the 4–momentum of the Higgs boson p^{β_i}. The standard coupling is split off explicitly such that $T_{\mu\nu}$ is normalized to be $g_{\mu\nu}$ in the SM. Moreover, with the assumption of massless leptons in the initial or final states, $T_{\mu\nu\beta_1\ldots\beta_\mathcal{J}}$ is transverse due to the conservation of the lepton currents, strongly constraining the form of the tensor.

Equivalent to the covariant description with the tensor $T_{\mu\nu\beta_1\ldots\beta_\mathcal{J}}$, the helicity formalism is the most convenient theoretical tool for defining observables which uniquely prove the scalar nature of the Higgs boson(s). In general, the basic helicity amplitudes [5] of the Higgs–strahlung and decay processes (1) for arbitrary H spin–$\mathcal{J}$, with the azimuthal angles set to zero, can be denoted by

$$\langle Z(\lambda_2)H(\lambda_H)|Z^*(\lambda_1)\rangle = \frac{g_W M_Z}{\cos\theta_W}\,\Gamma_{\lambda_2\lambda_H}\, d^1_{\lambda_1,\,\lambda_2-\lambda_H}(\theta); \tag{3}$$

$$\langle Z^{(*)}(\lambda_1)Z(\lambda_2)|H(\lambda_H)\rangle = \frac{g_W M_Z}{\cos\theta_W}\,\mathcal{T}_{\lambda_1\lambda_2}\, d^\mathcal{J}_{\lambda_H,\,\lambda_1-\lambda_2}(\theta) \tag{4}$$

respectively. No dependence of the reduced vertices $\Gamma_{\lambda_2\lambda_H}$ and $\mathcal{T}_{\lambda_1\lambda_2}$ on the Z^* spin component λ_1 and the H spin component λ_H is guaranteed by rotational invariance. The polar angle θ in Higgs–strahlung defines the momentum direction of the Z boson with helicity λ_2 with respect to the beam axis while the angle θ in the decay defines the polarization axis in the coordinate system with the momentum of the $Z^{(*)}$ boson with helicity λ_1 along the positive z–axis. The general helicity amplitudes and tensors for spins ≤ 2 can be found in Refs.[1] and [2], respectively.

The normality of the Higgs state, $n_H = (-1)^\mathcal{J}\,\mathcal{P}$, the product of the spin signature $(-1)^\mathcal{J}$ and the parity $\mathcal{P}$, connects the helicity amplitudes under parity transformations. If the vertex (2) are $\mathcal{P}$ invariant, equivalent to $\mathcal{CP}$ invariance in this specific case, the reduced vertices are related,

$$\Gamma_{\lambda_2\lambda_H} = n_H\,\Gamma_{-\lambda_2,-\lambda_H}; \qquad \mathcal{T}_{\lambda_1\lambda_2} = n_H\,\mathcal{T}_{-\lambda_1,-\lambda_2} \tag{5}$$

In addition, the helicity amplitudes $\mathcal{T}_{\lambda_1\lambda_2}$ are restricted further above the threshold of two real Z bosons by Bose symmetry; $\mathcal{T}_{\lambda_1\lambda_2} = (-1)^\mathcal{J}\,\mathcal{T}_{\lambda_2\lambda_1}$, independently of the parity of the H state.

In the SM the helicity amplitudes of Higgs–strahlung and the decay $H \to Z^{(*)}Z$ read

$$\Gamma_{00} = -\frac{E_Z}{M_Z}, \quad \Gamma_{10} = -1, \quad \Gamma_{01} = \Gamma_{11} = \Gamma_{12} = 0 \tag{6}$$

$$\mathcal{T}_{00} = \frac{M_H^2 - M_Z^2 - M_*^2}{2M_Z M_*}, \quad \mathcal{T}_{11} = -1, \quad \mathcal{T}_{10} = \mathcal{T}_{01} = \mathcal{T}_{1,-1} = 0 \tag{7}$$

with M_* the invariant mass of the Z^* boson, being M_Z if Z^* becomes real, and the SM Higgs boson carries even normality: $n_H = +1$.

The threshold behaviors in Higgs–strahlung and the decay $H \to Z^*Z$ with one virtual Z boson, which will play a key role in establishing the spin–parity of the Higgs boson(s), are determined by the leading β dependence of the helicity amplitudes with β, the Z three–momentum in the center of mass frame and in the H rest frame, respectively. The dependence can be worked out by counting the number of momenta in each term of the tensor $T_{\mu\nu\beta_1\ldots\beta_{\mathcal{J}}}$. Each momentum contracted with the Z–boson polarization vector or the H polarization tensor gives zero or one power of β. Furthermore, any momentum contracted with the lepton current also gives rise to one power of β due to the transversality of the current. The overall β dependence can be derived from the squared β dependence of the helicity amplitude multiplied by a single factor β from the phase space.

Based on the general description of the HZZ vertex, we will demonstrate in the consecutive sections the unique spin–parity determination of the Higgs boson(s) for a $\mathcal{CP}$ invariant theory, for even and odd normality Higgs bosons in Higgs–strahlung $e^+e^- \to ZH$ and the decay process $H \to ZZ$. The analyses can straightforwardly be extended to mixed parity assignment in $\mathcal{CP}$ noninvariant theories.

3 Higgs–strahlung

3.1 Characteristic observables

The total cross section for Higgs–strahlung in a $\mathcal{CP}$ invariant theory is given in terms of the reduced vertex $\Gamma_{\lambda_2\lambda_H}$ by

$$\sigma(e^+e^- \to ZH) \sim \beta \left[|\Gamma_{00}|^2 + 2(|\Gamma_{11}|^2 + |\Gamma_{01}|^2 + |\Gamma_{10}|^2 + |\Gamma_{12}|^2)\right] \tag{8}$$

In the SM with the helicity amplitudes (6) the cross section rise linearly with β, i.e., steeply with the center of mass energy above the threshold: $\sigma_H \sim \beta \sim [s-(M_H+M_Z)^2]^{1/2}$. This steep rise is characteristic of the production of a scalar particle in conjunction with the Z boson (with only two exceptions as discussed below).

The distribution of the polar angle θ can in general be written

$$\frac{d\sigma}{d\cos\theta} \sim \sin^2\theta \left(|\Gamma_{00}|^2 + 2|\Gamma_{11}|^2\right) + (1+\cos^2\theta)\left(|\Gamma_{01}|^2 + |\Gamma_{10}|^2 + |\Gamma_{12}|^2\right) \tag{9}$$

The SM distribution, $d\sigma_H/d\cos\theta \sim \beta^2\sin^2\theta + 8M_Z^2/s$, is isotropic near the threshold and it approaches $\sim \sin^2\theta$ at high energies, congruent with the equivalence theorem.

Other independent information on the helicity of the Z state is encoded in the final–state fermion distribution in the decay $Z \to f\bar{f}$. Denoting the fermion polar angle in the Z rest frame with respect to the Z flight direction in the laboratory frame by θ_*, the general form of the double differential distribution in θ and θ_* reads

$$\frac{d\,\sigma}{d\cos\theta d\cos\theta_*} \sim \sin^2\theta\sin^2\theta_*|\Gamma_{00}|^2$$

$$+\frac{1}{2}\left[(1+\cos^2\theta)(+\cos^2\theta_*)\left(|\Gamma_{10}|^2+|\Gamma_{12}|^2\right)+4\eta_e\eta_f\cos\theta\cos\theta_*\left(|\Gamma_{10}|^2-|\Gamma_{12}|^2\right)\right]$$
$$+\sin^2\theta(1+\cos^2\theta_*)|\Gamma_{11}|^2+(1+\cos^2\theta)\sin^2\theta_*|\Gamma_{01}|^2 \qquad (10)$$

where $\eta_f = 2v_f a_f/(v_f^2+a_f^2)$ with the electroweak charges $v_f = 2I_{3f} - 4e_f\sin^2\theta_W$ and $a_f = 2I_{3f}$ of the fermion f. The SM correlated angular distribution has *no* $\sin^2\theta(1+\cos^2\theta_*)$ *and no* $(1+\cos^2\theta)\sin^2\theta_*$ *terms* and approaches $\sim \sin^2\theta\sin^2\theta_*$ for high energies, reflecting the longitudinal Z polarization in the asymptotic limit.

3.2 Selection rules

States of *even normality* $\mathcal{J}^P = 1^-, 2^+, 3^-, \ldots$ can be excluded by measuring the threshold behaviour of the excitation curve and the angular correlations:

Spin 1: All helicity amplitudes vanish near threshold linearly in β, so the excitation curve rises $\sim \beta^3$, distinct from the SM.

Spin 2: The most general spin–2 tensor contains a term with no momentum dependence, resulting in non–zero helicity amplitudes at threshold. However, Γ_{01} and Γ_{11} are non–zero in this case, leading to non–trivial $(1+\cos^2\theta)\sin^2\theta_*$ and $\sin^2\theta(1+\cos^2\theta_*)$ correlations in Eq.(10) which are absent in the SM. Therefore, if the excitation curve rises linearly, not observing these correlations rules out the spin–2 assignment to the H state. If the term in the spin–2 case vanishes, the excitation curve rises $\sim \beta^5$ near threshold.

Spin≥ 3: Above spin–2 the number of independent helicity amplitudes does not increase any more [5] and the most general spin–$\mathcal{J}$ tensor $T_{\mu\nu\beta_1\ldots\beta_\mathcal{J}}$ is a direct product of a tensor $T^{(2)}_{\mu\nu\beta_i\beta_j}$ isomorphic with the spin–2 tensor and a symmetric tensor built up by the momentum vectors $q^{\beta_k} = (p_Z + p_H)^{\beta_k}$ as required by the properties of the spin–$\mathcal{J}$ wave–function $H^{\beta_1\ldots\beta_\mathcal{J}}$. Contracted with the wave–function, the extra $\mathcal{J}$–2 momentum give rise to a leading power $\beta^{\mathcal{J}-2}$ in the helicity amplitudes. The cross section therefore rises near threshold $\sim \beta^{2\mathcal{J}-3}$, i.e., with a power ≥ 3, in contrast to the single power of the SM.

It is quite easy to rule out particles of *odd normality*: $\mathcal{J}^\mathcal{P} = 0^-, 1^+, 2^-, \ldots$. Since the helicity amplitude Γ_{00} must vanish by the $\mathcal{CP}$ relation (5), the observation of a non–zero $\sin^2\theta\sin^2\theta_*$ correlation, as predicted by the SM, eliminates all odd normality states. We find a similar picture to the even normality case, where the excitation curve only presents a linear rise for a particle of spin–1. The generalization to higher spins ≥ 3 follows exactly as before, resulting in an excitation curve $\sim \beta^{2\mathcal{J}-1}$, i.e., with a power ≥ 5, at threshold.

4 Higgs decays to Z pairs

4.1 Characteristic observables

The key observables for measuring the spin–parity of the Higgs boson(s) are the invariant mass (M_*) spectrum of the off–shell Z boson in the decay $H \to Z^*Z$ and the angular distributions of the final–state fermions in the decays $Z^{(*)} \to f\bar{f}$, encoding the helicities of the $Z^{(*)}$ states. The combined polar and azimuthal angular distributions are presented in the Appendix of Ref.[2].

Below the threshold of two real Z bosons, the invariant mass (M_*) spectrum of the off–shell Z boson is maximal close to the kinematical limit corresponding to zero momentum of the off– and on–shell Z bosons and it decreases linearly with β, i.e., steeply with the invariant mass just below the threshold: $d\Gamma_H/dM_*^2 \sim \beta \sim [(M_H - M_Z)^2 - M_*^2]^{1/2}$. This steep decrease is characteristic of the decay of a scalar particle into two vector bosons (with only two exceptions as discussed below).

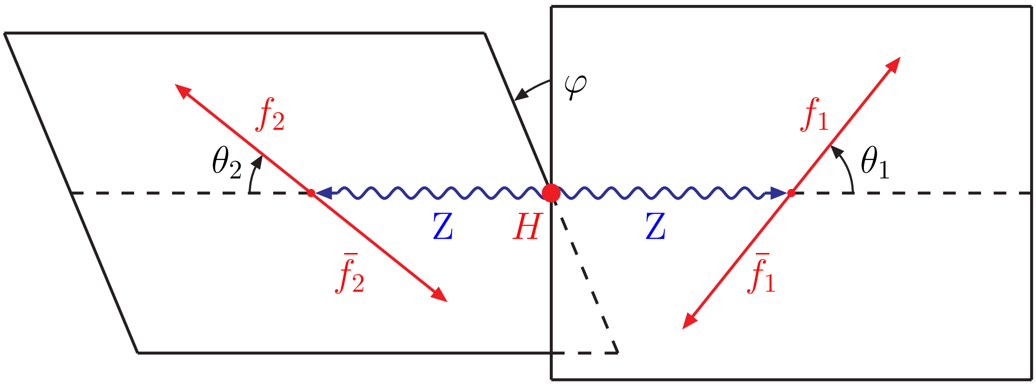

Figure 1: *The definition of the polar angles θ_i $(i = 1, 2)$ and the azimuthal angle φ for the sequential decay $H \to Z^{(*)}Z \to (f_1\bar{f}_1)(f_2\bar{f}_2)$ in the rest frame of the Higgs particle.*

Polar and azimuthal angular distributions give independent access to spin and parity of the Higgs boson. Denoting the polar angles of the fermions f_1, f_2 in the rest frames of the Z bosons by θ_1 and θ_2, and the azimuthal angle between the planes of the fermion pairs by φ [see Fig. 1], the polar–angle distributions for a $\mathcal{CP}$ invariant theory can be written as

$$\begin{aligned}\frac{d\Gamma}{d\cos\theta_1 d\cos\theta_2} \sim\ & \sin^2\theta_1 \sin^2\theta_2\, |\mathcal{T}_{00}|^2 + \frac{1}{2}(1+\cos^2\theta_1)(1+\cos^2\theta_2)\left[|\mathcal{T}_{11}|^2 + |\mathcal{T}_{1,-1}|^2\right] \\ & + (1+\cos^2\theta_1)\sin^2\theta_2\, |\mathcal{T}_{10}|^2 + \sin^2\theta_1\,(1+\cos^2\theta_2)\, |\mathcal{T}_{01}|^2 \\ & + 2\,\eta_1\eta_2 \cos\theta_1\cos\theta_2 \left[|\mathcal{T}_{11}|^2 - |\mathcal{T}_{1,-1}|^2\right] \qquad (11)\end{aligned}$$

while the general azimuthal angular distribution reads

$$\begin{aligned}\frac{d\Gamma}{d\varphi} \sim\ & |T_{11}|^2 + |T_{10}|^2 + |T_{1,-1}|^2 + |T_{01}|^2 + |T_{00}|^2/2 \\ & + \eta_1\eta_2\left(\frac{3\pi}{8}\right)^2 \Re\mathrm{e}(T_{11}T_{00}^* + T_{10}T_{0,-1}^*)\cos\varphi + \frac{1}{4}\Re\mathrm{e}(T_{11}T_{-1,-1}^*)\cos 2\varphi \qquad (12)\end{aligned}$$

The SM polar–angle distributions with the helicity amplitudes (7) approach $\sim \sin^2\theta_1 \sin^2\theta_2$ for large Higgs masses, reflecting the longitudinal Z polarization [4]. Also any φ dependence disappears in this limit.

4.2 Selection rules

4.2.1 Heavy Higgs bosons

Since the helicity amplitude $\mathcal{T}_{00}$ must vanish by the $\mathcal{CP}$ relation (5) for *odd normality*, observing a non-zero $\sim \sin^2\theta_1 \sin^2\theta_2$ correlation as predicted by the SM, eliminates all odd–normality states.

In the chain of *even-normality* states, the odd-spin states 1^-, 3^-, …, can easily be excluded by observing the $\sin^2\theta_1 \sin^2\theta_2$ correlation induced by $\mathcal{T}_{00}$ in the Standard Model, but forbidden by Bose symmetry for even-normality odd-spin states.

In general, the vertex for the higher even–$\mathcal{J}$ Higgs state will lead to four–fermion angular correlations different from those for the spin–0 case. If $T_{\mu\nu\beta_1\ldots\beta_\mathcal{J}} = \left[T^{\mathcal{J}=0}_{\mu\nu}\right] k_{\beta_1}\ldots k_{\beta_\mathcal{J}}$ [with $k = k_1 - k_2$], however, the unpolarized higher even–$\mathcal{J}$ state generates the same angular correlations of the Z decay products as the spin–0 state. Thus, from final-state distributions alone, a model–independent spin–parity analysis cannot be carried out. However, special production mechanisms such as $gg \to H$ at LHC [6] and $\gamma\gamma \to H$ in the Compton mode of linear colliders [7] can be successfully exploited to close the gap. Assuming the HZZ coupling to be of the special form, the polar–angle distribution for the process $gg/\gamma\gamma \to H \to ZZ$ is given by the differential cross section

$$\frac{d\sigma}{d\cos\theta}\,[gg/\gamma\gamma \to H \to ZZ] \sim |a_1|^2 \left[P^0_\mathcal{J}(\cos\theta)\right]^2 + 12|a_2|^2 \left[P^2_\mathcal{J}(\cos\theta)\right]^2 \qquad (13)$$

where a_1, a_2 are two independent form factors describing the general spin–$\mathcal{J}$ tensor for the ggH or $\gamma\gamma H$ coupling [2] and θ is the polar angle between the momenta of a gluon/photon and a Z boson in the $gg/\gamma\gamma$ center–of–mass frame. The associated Legendre functions $\mathcal{P}^2_\mathcal{J}$ and $\mathcal{P}^0_\mathcal{J}$ have non-trivial $\cos\theta$ dependence except for $\mathcal{J} = 0$. Thus, the zero–spin of the Higgs boson can be checked through the lack of the polar (and azimuthal) angle correlations between the initial state and final state particles in the combined process, $gg/\gamma\gamma \to H \to ZZ$.

4.2.2 Intermediate Higgs-mass range

The spin–parity analysis in Higgs decays to a pair of virtual and real Z bosons, $H \to Z^*Z$, runs parallel in all elements to the same task in Higgs-strahlung at e^+e^- colliders – just requiring the crossing of the virtual Z-boson line from the initial to the final state.

For the same arguments as before, the states of *odd normality* can be excluded if a non-zero $\sim \sin^2\theta_1 \sin^2\theta_2$ correlation has been established experimentally. Equivalently, the high power suppression of the virtual mass distributions near the threshold rules out all spin ≥ 2 states; the state $\mathcal{J} = 1$ can be eliminated by non-observation of $\sim (1+\cos^2\theta_1)\sin^2\theta_2$ and $\sin^2\theta_1 (1+\cos^2\theta_2)$ correlations.

Below the threshold of two real Z bosons, the states of *even normality* can be excluded by measuring the threshold behaviour of the invariant mass spectrum and the angular correlations.

Spin 1: Every helicity amplitude vanishes near threshold linearly in β, so the invariant mass spectrum decreases $\sim \beta^3$, distinct from the SM.

Spin 2: A momentum–independent term, $T^{\mu\nu\beta_1\beta_2} \sim g^{\mu\beta_1}g^{\nu\beta_2} + g^{\mu\beta_2}g^{\nu\beta_1}$, resulting in non–zero helicity amplitudes at threshold, generates the helicity amplitudes $\mathcal{T}_{10}$ and $\mathcal{T}_{01}$, leading to non-trivial $(1+\cos^2\theta_1)\sin^2\theta_2$ and $\sin^2\theta_1(1+\cos^2\theta_2)$ correlations absent in the SM. Therefore, if the invariant mass spectrum decreases linearly not observing these polar–angle correlations rules out the spin–2 assignment to the state. Without this peculiar term in the spin-2 case, the spectrum falls off $\sim \beta^5$ near threshold.

Spin $\geq$ 3: Contracted with the wave–function, the extra $\mathcal{J}-2$ momenta in the general spin–$\mathcal{J}$ tensor $T_{\mu\nu\beta_1\ldots\beta_\mathcal{J}}$ give rise to a leading power $\beta^{\mathcal{J}-2}$ in the helicity amplitudes. The invariant mass spectrum therefore decreases near threshold $\sim \beta^{2\mathcal{J}-3}$, *i.e.* with a power ≥ 3, in contrast to the single power of the SM.

5 Conclusions

The spin–parity analyses described above can be briefed in a few characteristic points.

In Higgs–strahlung at e^+e^- linear colliders the rise of the excitation curve near the threshold and the angular distributions render the spin–parity analysis of the Higgs boson unambiguous.

Complementary to the spin–parity measurements in Higgs–strahlung, Higgs decays to Z bosons can provide us with a clear picture of these external quantum numbers for Higgs masses above the ZZ threshold, if auxiliary angular distributions are included that are generated in specific production mechanism such as gluon fusion at the LHC and $\gamma\gamma$ fusion at linear colliders. Below the mass range for on-shell ZZ decays, threshold analyses combined with angular correlations in Z^*Z decays [with one of the electroweak bosons, Z^*, being virtual] can be exploited in analogy to Higgs-strahlung.

The rules can be supplemented by observations specific to two cases. By observing non–zero $H\gamma\gamma$ and Hgg couplings, the $\mathcal{J}=1$ assignment can elegantly be ruled out by Yang's theorem in particular, and for all odd spins in general.

The above formalisms can be generalized easily to rule out mixed normality states with spin ≥ 1. For a Higgs boson of mixed normality we cannot use Eq.(5) anymore to derive the simple form of the differential distributions. The analysis for identifying the spin of the Higgs particle, however, proceeds exactly as before in the fixed normality case, since the most general vertex will be the sum of the even and odd normality tensors.

Before closing, we note that experimental simulations for the spin–parity determinations have been performed for Higgs decays to Z boson pairs at the LHC in Refs.[2, 8] and for Higgs–strahlung at linear colliders in Ref.[9].

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