# Spin and orbital structure of $f\bar{f}$ in scalar and pseudoscalar decays
# Спиновая и орбитальная структура $f\bar{f}$ в распадах скаляра и псевдоскаляра

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Мы вычисляем волновую функцию в координатном представлении для пары фермион–антифермион, рождающейся в двухчастичных распадах спин-нулевой частицы. Исходя из амплитуды $S$-матрицы, мы получаем волновую функцию в импульсном представлении и её преобразование Фурье. Для скалярной $(0^+)$ связи получается чистая парциальная волна с $L = 1$, при этом спиновые множители фиксируют определённые проекции орбитального момента $m$ для конкретных ориентаций спинов $f$ и $\bar{f}$. Для псевдоскалярной $(0^-)$ связи результатом является чистая $L = 0$ волна с обычной синглетной спиновой структурой. Мы обсуждаем, почему противоположно направленные импульсы конечных частиц не противоречат ненулевому орбитальному моменту: $L$ является свойством относительной координаты. Также отмечено, почему фазы орбитального момента, как правило, невидимы в стандартных измерениях физики высоких энергий, основанных на проекции на собственные состояния импульса, и указаны спин-разрешённые угловые корреляции как наблюдаемые проявления этих фаз.

We construct the coordinate-space wave function of a fermion–antifermion pair produced in two-body decays of a spin-zero parent. Starting from the $S$-matrix amplitude we derive a momentum-space wave function and its Fourier transform. For a scalar $(0^+)$ coupling we obtain a pure $L = 1$ partial wave whose spin factors select definite orbital projections $m$ for specific spin orientations of $f$ and $\bar{f}$. For a pseudoscalar $(0^-)$ coupling the result is a pure $L = 0$ wave with the usual singlet spin structure. We show explicitly why the back-to-back momenta of the final particles do not contradict nonzero orbital angular momentum: $L$ is a property of the relative coordinate. We also discuss why OAM phases are generally invisible in standard HEP measurements that project onto momentum eigenstates, and point to spin-resolved angular correlations as the observable remnants.

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## 1. Introduction

Two-body decays of spinless parents admit both $L = 0$ and $L \neq 0$ final states by symmetry. Momentum conservation, however, enforces back-to-back kinematics in the parent rest frame, which can obscure how nonzero orbital angular momentum (OAM) enters the problem.

Starting from the $S$-matrix element for $0^{\pm} \to f\bar{f}$ we construct the properly normalized momentum-space amplitude and its coordinate-space wave function $\psi(\mathbf{r}; \mathbf{s}_1, \mathbf{s}_2)$, with $\mathbf{r}$ the relative coordinate and $\mathbf{s}_{1,2}$ the spin quantization axes. For a scalar ($0^+$) coupling,

$$\psi(\mathbf{r}) \propto j_1(pr)\, \chi^{\dagger}(\mathbf{s}_1)\, (\boldsymbol{\sigma}\cdot\hat{\boldsymbol{r}})\, \eta(\mathbf{s}_2),$$

a pure $L = 1$ (P-wave). The decomposition of $\boldsymbol{\sigma}\cdot\hat{\boldsymbol{r}}$ into spherical tensors shows that fixing the spin projections selects a definite $Y_{1m}$ and therefore a definite OAM projection $m$.

For a pseudoscalar ($0^-$) coupling,

$$\psi(\mathbf{r}) \propto j_0(pr)\, \chi^{\dagger}(\mathbf{s}_1)\eta(\mathbf{s}_2),$$

a pure $L = 0$ (S-wave). Here $j_0$ and $j_1$ are Bessel functions of the first kind of order zero and one, respectively, while $\chi$ and $\eta$ are Pauli spinors. Since standard HEP observables integrate over $\mathbf{r}$ and project onto momentum eigenstates, global OAM phases cancel; the measurable quantities are the resulting spin–angle correlations. These correlations, derived below in closed form, provide a direct probe of the scalar, pseudoscalar, or mixed nature of the coupling.

Spin–orbital correlations are known to provide direct sensitivity to the CP structure of Yukawa interactions. In the $H \to \tau^+\tau^-$ channel, for example, CMS used the relative orientation of the tau decay planes to constrain the CP phase of the Yukawa coupling [1]. Here we analyze the same class of correlations at the level of the coordinate-space wave function for $0^{\pm} \to f\bar{f}$, where the origin of the CP-sensitive structures becomes explicit.

The paper is organized as follows. Sections 2–4 present the scalar, pseudoscalar, and mixed cases, respectively. In Sec. 5 we discuss the results and briefly review the related literature. Conclusions are given in Sec. 6.

Natural units $\hbar = c = 1$ are used throughout.

## 2. Scalar case: coordinate-space wave function

The single-particle states are relativistic normalized:

$$\langle \mathbf{p}'|\mathbf{p}\rangle = (2\pi)^3 2E_{\mathbf{p}}\delta^3(\mathbf{p} - \mathbf{p}').$$

We consider the decay of a scalar Higgs boson, $H \to f\bar{f}$, into a fermion $f$ and antifermion $\bar{f}$ with spin 3-vectors $\mathbf{s}_1$ and $\mathbf{s}_2$. The final state can be decomposed in momentum space as

$$|f\bar{f}\rangle = \int \frac{d\mathbf{p}_1}{(2\pi)^3}\frac{d\mathbf{p}_2}{(2\pi)^3}\frac{\widetilde{\psi}_{r_1 r_2}(\mathbf{p}_1, \mathbf{p}_2)}{\sqrt{2E_{\mathbf{p}_1}2E_{\mathbf{p}_2}}}|\mathbf{p}_1, r_1; \mathbf{p}_2, r_2\rangle, \tag{1}$$

where $r_i$ is the spin projection on a fixed $z$-axis in configuration space and

$$|\mathbf{p}_1, r_1; \mathbf{p}_2, r_2\rangle = \sqrt{2E_{\mathbf{p}_1} 2E_{\mathbf{p}_2}}\, a^\dagger_{r_1}(\mathbf{p}_1)\, b^\dagger_{r_2}(\mathbf{p}_2)\,|0\rangle$$

with $a^\dagger, b^\dagger$ the creation operators for $f$ and $\bar{f}$. The (unnormalized) coordinate-space wave function is the Fourier transform of $\widetilde{\psi}$,

$$\psi_{r_1 r_2}(\mathbf{x}_1, \mathbf{x}_2) = \int \frac{d\mathbf{p}_1}{(2\pi)^3}\frac{d\mathbf{p}_2}{(2\pi)^3}\widetilde{\psi}_{r_1 r_2}(\mathbf{p}_1, \mathbf{p}_2) e^{i(\mathbf{p}_1\cdot\mathbf{x}_1 + \mathbf{p}_2\cdot\mathbf{x}_2)}. \tag{2}$$

Projecting the quantum state $S|H\rangle$ onto the two-particle Fock state gives

$$\widetilde{\psi}_{r_1 r_2}(\mathbf{p}_1, \mathbf{p}_2) = \frac{\langle \mathbf{p}_1, r_1; \mathbf{p}_2, r_2|S|H\rangle}{\sqrt{2E_{\mathbf{p}_1}\, 2E_{\mathbf{p}_2}}} = (2\pi)^4\delta^4(P - p_1 - p_2)\,\frac{i\mathcal{M}_{r_1 r_2}(\mathbf{p}_1, \mathbf{p}_2)}{\sqrt{2E_{\mathbf{p}_1}\, 2E_{\mathbf{p}_2}}}\,, \tag{3}$$

where $P$ is the four-momentum of the Higgs boson. At leading order in the Standard Model,

$$i\mathcal{M}_{r_1 r_2}(\mathbf{p}_1, \mathbf{p}_2) = i\,\frac{m_f}{v}\,\overline{u}_{r_1}(\mathbf{p}_1)\, v_{r_2}(\mathbf{p}_2)\,, \tag{4}$$

with $m_f$ the fermion mass and $v$ the Higgs vacuum expectation value.

Using Eqs. (2)–(4) and going to the Higgs rest frame, we obtain

$$\psi_{r_1 r_2}(\mathbf{r}) = i\,\frac{m_f\, p}{16\pi^2 v}\int d\Omega_{\mathbf{n}}\; e^{ip\,\mathbf{n}\cdot\mathbf{r}}\,\overline{u}_{r_1}(p\mathbf{n})\, v_{r_2}(-p\mathbf{n})\,, \tag{5}$$

where $\mathbf{r} = \mathbf{x}_1 - \mathbf{x}_2$ is the relative coordinate, $E_p = m_H/2$, $p = \sqrt{E_p^2 - m_f^2}$, and the integral is taken over the unit sphere.

To allow arbitrary spin quantization axes $\mathbf{s}_1, \mathbf{s}_2$ we rotate the $z$-basis spinors by the SU(2) matrices

$$D^{(1/2)}(\mathbf{s}) = e^{-i\phi\sigma_z/2}\, e^{-i\theta\sigma_y/2} = \begin{pmatrix} \cos\frac{\theta}{2} & -e^{-i\phi}\sin\frac{\theta}{2} \\ e^{i\phi}\sin\frac{\theta}{2} & \cos\frac{\theta}{2} \end{pmatrix}, \tag{6}$$

where $\mathbf{s} = (\sin\theta\cos\phi, \sin\theta\sin\phi, \cos\theta)$.

The wave function with definite spin projections $\rho_{1,2} = \pm 1$ along $\mathbf{s}_{1,2}$ can be written as

$$\psi_{\rho_1\rho_2}(\mathbf{r}; \mathbf{s}_1, \mathbf{s}_2) = \sum_{r_1, r_2 = \pm} D^{(1/2)}_{r_1\rho_1}(\mathbf{s}_1)\, D^{(1/2)}_{r_2\rho_2}(\mathbf{s}_2)\,\psi_{r_1 r_2}(\mathbf{r}). \tag{7}$$

Equivalently, we can absorb the rotations directly into the Dirac spinors,

$$\begin{aligned} u_\rho(p\mathbf{n}, \mathbf{s}) &= \begin{pmatrix} \sqrt{E_p + m_f}\,\chi_\rho(\mathbf{s}) \\ \sqrt{E_p - m_f}\,(\boldsymbol{\sigma}\cdot\mathbf{n})\,\chi_\rho(\mathbf{s}) \end{pmatrix}, \\ v_\rho(-p\mathbf{n}, \mathbf{s}) &= \begin{pmatrix} -\sqrt{E_p - m_f}\,(\boldsymbol{\sigma}\cdot\mathbf{n})\,\eta_\rho(\mathbf{s}) \\ \sqrt{E_p + m_f}\,\eta_\rho(\mathbf{s}) \end{pmatrix}, \end{aligned} \tag{8}$$

where $\chi_\rho(\mathbf{s})$ and $\eta_\rho(\mathbf{s}) = i\sigma^2\chi^*_\rho(\mathbf{s})$ are the Pauli spinors for $f$ and $\bar{f}$.

Inserting these expressions into Eq. (5) one finds the coordinate-space wave function

$$\psi_{\rho_1\rho_2}(\mathbf{r};\mathbf{s}_1,\mathbf{s}_2) = \frac{m_f\,p^2}{2\pi v}\,j_1(pr)\,\chi^\dagger_{\rho_1}(\mathbf{s}_1)\,(\boldsymbol{\sigma}\cdot\hat{\boldsymbol{r}})\,\eta_{\rho_2}(\mathbf{s}_2). \tag{9}$$

To obtain $|\psi|^2$ we use the polarization projectors

$$\chi_\lambda(\mathbf{s})\,\chi^\dagger_\lambda(\mathbf{s}) = \frac{1}{2}\Big(1+\lambda\,\mathbf{s}\cdot\boldsymbol{\sigma}\Big), \qquad \eta_\lambda(\mathbf{s})\,\eta^\dagger_\lambda(\mathbf{s}) = \frac{1}{2}\Big(1-\lambda\,\mathbf{s}\cdot\boldsymbol{\sigma}\Big) \tag{10}$$

$$|\psi_{\rho_1\rho_2}(\mathbf{r};\mathbf{s}_1,\mathbf{s}_2)|^2 = \Big(\frac{m_f\,p^2}{2\pi v}\Big)^2 j_1^2(pr)\;\frac{1}{2}\Big[\,1\;+\;\rho_1\rho_2\big(\mathbf{s}_1\cdot\mathbf{s}_2-2(\mathbf{s}_1\cdot\hat{\boldsymbol{r}})(\mathbf{s}_2\cdot\hat{\boldsymbol{r}})\big)\Big]. \tag{11}$$

where $\hat{\boldsymbol{r}} = \mathbf{r}/r$.

The scalar wave function in (9) contains the operator $\boldsymbol{\sigma}\cdot\hat{\mathbf{r}}$, so the spatial part is a pure $L = 1$ wave.

Two geometries illustrate the mapping to definite OAM projections: (i) $\hat{\mathbf{r}} \parallel \hat{\mathbf{s}}$: only anti-parallel spin outcomes survive, selecting a pure $m = 0$ component; (ii) $\hat{\mathbf{r}} \perp \hat{\mathbf{s}}$: only parallel outcomes survive, selecting the $m = \pm 1$ sector (sign fixed by $\rho_{1,2}$).

Thus spin eigenstates along a common axis project the relative motion onto a single spherical–harmonic eigenstate within the $L = 1$ multiplet.

## 3. Pseudoscalar ($0^-$) case

For a pseudoscalar coupling the amplitude reads

$$i\mathcal{M}_{\rho_1\rho_2} = i\,\frac{m_f}{v}\,\overline{u}_{\rho_1}(p\mathbf{n},\mathbf{s}_1)\,\gamma^5\,v_{\rho_2}(-p\mathbf{n},\mathbf{s}_2), \tag{12}$$

and with the standard Dirac spinors in the Higgs rest frame one finds

$$\overline{u}_{\rho_1}(p\mathbf{n},\mathbf{s}_1)\,\gamma^5\,v_{\rho_2}(-p\mathbf{n},\mathbf{s}_2) = 2E_p\,\chi^\dagger_{\rho_1}(\mathbf{s}_1)\,\eta_{\rho_2}(\mathbf{s}_2), \tag{13}$$

independent of $\mathbf{n}$. The coordinate-space wave function is therefore

$$\psi_{\rho_1\rho_2}(\mathbf{r};\mathbf{s}_1,\mathbf{s}_2) = \frac{m_f E_p}{2\pi v}\,j_0(pr)\;\chi^\dagger_{\rho_1}(\mathbf{s}_1)\,\eta_{\rho_2}(\mathbf{s}_2), \tag{14}$$

a pure $L = 0$ partial wave. Its modulus squared reads

$$|\psi_{\rho_1\rho_2}|^2 = \Big(\frac{m_f E_p}{2\pi v}\Big)^2 j_0^2(pr)\;\frac{1}{2}\Big[1-\rho_1\rho_2\,\mathbf{s}_1\cdot\mathbf{s}_2\Big]. \tag{15}$$

Thus the pseudoscalar populates the S-wave singlet: for $\mathbf{s}_1 = \mathbf{s}_2$ only anti-parallel outcomes survive, and there is no dependence on $\hat{\boldsymbol{r}}$.

## 4. Mixed scalar–pseudoscalar coupling

For a general complex coupling $\varepsilon_1 H\bar{f}f + i\varepsilon_2 H\bar{f}\gamma^5 f$ the coordinate-space wave function is a superposition

$$\psi = \varepsilon_1\,\psi_S + \varepsilon_2\,\psi_P, \tag{16}$$

where $\psi_S$ and $\psi_P$ are the scalar and pseudoscalar wave functions discussed above. The probability density decomposes as

$$|\psi|^2 = |\varepsilon_1|^2|\psi_S|^2 + |\varepsilon_2|^2|\psi_P|^2 + 2\,\Re\big(\varepsilon_1\varepsilon_2^*\,\psi_S\psi_P^*\big). \tag{17}$$

The first two terms reproduce Eqs. (11) and the pseudoscalar result. The interference term is odd under naive time reversal and is governed by a triple product of the form

$$\psi_S\psi_P^* \;\propto\; j_0(pr)\,j_1(pr)\,\rho_1\rho_2\,\mathbf{s}_1\cdot(\hat{\boldsymbol{r}}\times\mathbf{s}_2), \tag{18}$$

so that

$$|\psi|^2 \supset \Re(\varepsilon_1\varepsilon_2^*)\,j_0(pr)\,j_1(pr)\,\rho_1\rho_2\,\mathbf{s}_1\cdot(\hat{\boldsymbol{r}}\times\mathbf{s}_2), \tag{19}$$

up to an overall prefactor common to all terms. This structure provides a natural CP-sensitive observable in spin–angle correlations.

## 5. Discussion

The determination of the $CP$ properties of the Higgs boson has been a central topic since its discovery in 2012 [2, 3]. Early analyses of $H \to ZZ$ and $H \to WW$ decays excluded a pure pseudoscalar hypothesis [2, 3]. However, these channels probe the Higgs coupling to gauge bosons and are only indirectly sensitive to possible $CP$-odd admixtures in Yukawa interactions. Direct access to the $CP$ structure of the fermionic coupling is provided by $H \to \tau^+\tau^-$, where the parity-violating weak decays of the $\tau$ leptons act as efficient spin analyzers.

A substantial phenomenological program has been developed to extract the scalar–pseudoscalar mixing angle $\phi_\tau$ (or equivalently $\alpha^{H\tau\tau}$) from angular correlations in $H \to \tau\tau$ decays at hadron colliders [4–13]. Two complementary strategies are commonly employed. The impact-parameter method applies to all $\tau$ decay modes with charged final states [4–7]. The $\rho$-decay plane method exploits the enhanced sensitivity of the channel in which both $\tau$ leptons decay via intermediate $\rho$ mesons [9–13]. In both approaches the essential observable is a $CP$-odd angular correlation constructed from the visible decay products.

A comprehensive phenomenological analysis was performed in Ref. [8], where the authors defined the signed angle $\phi^*_{CP}$ between the two $\tau$ decay planes and studied its sensitivity to the mixing angle. Including Drell–Yan background and detector effects, they estimated that $\phi_\tau$ could be measured

with an uncertainty of $\Delta\phi_\tau \approx 15^\circ$ ($9^\circ$) at the LHC with $150\,\mathrm{fb}^{-1}$ ($500\,\mathrm{fb}^{-1}$), improving to about $4^\circ$ at $3\,\mathrm{ab}^{-1}$. They also demonstrated that the Drell–Yan background produces an approximately flat $\phi^*_{CP}$ distribution, since it does not contain the scalar–pseudoscalar interference structure characteristic of a mixed Yukawa coupling.

These developments culminated in the first experimental measurement by the CMS collaboration [14]. Using the full Run 2 data set at $\sqrt{s} = 13\,\mathrm{TeV}$ corresponding to $137\,\mathrm{fb}^{-1}$, CMS reconstructed the acoplanarity angle $\phi_{CP}$ in several decay channels ($\tau_\mu\tau_h$, $\tau_e\tau_h$, and $\tau_h\tau_h$) and extracted the effective mixing angle $\alpha^{H\tau\tau} = -1 \pm 19^\circ$, consistent with the Standard Model prediction. A pure $CP$-odd coupling was excluded at the $3\sigma$ level.

The mixed scalar–pseudoscalar wave function derived in this work provides a direct theoretical origin of such $CP$-sensitive observables. In Eq. (19) we obtained an interference contribution to the probability density of the form

$$\propto\ \hat{\boldsymbol{r}} \cdot \big(\mathbf{s}_1 \times \mathbf{s}_2\big),$$

where $\hat{\boldsymbol{r}}$ is the unit vector along the relative coordinate. In the Higgs rest frame, $\hat{\boldsymbol{r}}$ is aligned with the fermion momentum direction, so this structure corresponds to a spin–momentum triple product. It is odd under $CP$ and arises solely from the interference between the scalar ($L = 1$) and pseudoscalar ($L = 0$) components of the wave function.

The experimentally used observables are momentum-space realizations of this structure. In particular:

1. The triple-product correlation constructed in Ref. [8],

$$\mathcal{O}^* = \hat{\mathbf{q}}^{*-} \cdot \big(\hat{\mathbf{n}}^+_\perp \times \hat{\mathbf{n}}^-_\perp\big),$$

   is an operational proxy for the fundamental correlation $\hat{\boldsymbol{r}} \cdot (\mathbf{s}_1 \times \mathbf{s}_2)$, with the $\tau$ decay products encoding the parent spin directions.

2. The signed acoplanarity angle $\phi^*_{CP}$ (or $\phi_{CP}$ in the CMS analysis [14]) reproduces the characteristic $2\phi_\tau$ phase dependence generated by the scalar–pseudoscalar interference term.

3. The near-flatness of the Drell–Yan background in $\phi^*_{CP}$ follows from the absence of a $CP$-odd scalar–pseudoscalar interference contribution in vector-current mediated processes.

Thus, the coordinate-space formalism developed here makes explicit the quantum-mechanical origin of the $CP$-sensitive angular correlations employed in collider analyses. The triple-product structure appears already at the level of the two-particle wave function and reflects the interference between distinct orbital components of the final state.

## 6. Conclusions

We derived the coordinate-space wave functions for the decays $0^{\pm} \to f\bar{f}$ with arbitrary spin quantization axes. The scalar coupling produces a pure $L = 1$ state whose OAM projection $m$ is fixed by the spin configuration, while the pseudoscalar coupling yields a pure $L = 0$ singlet. These analytic results make the spin–angle correlations fully transparent and provide a direct link between the orbital structure of the final state and the CP properties of the Yukawa interaction. The framework also applies to mixed scalar–pseudoscalar couplings, where the interference term generates a CP-sensitive triple product.

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