Determining the CP Property of $ht\bar{t}$ Coupling via a Novel Jet Substructure Observable

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Determining the CP property of the Higgs boson is important for a precision test of the Standard Model as well as for the search for new physics. We propose a novel jet substructure observable based on the azimuthal anisotropy in a linearly polarized gluon jet that is produced in association with a Higgs boson at hadron colliders, and demonstrate that it provides a new CP-odd observable for determining the CP property of the Higgs-top interaction. We introduce a factorization formalism to define a polarized gluon jet function with the insertion of an infrared-safe azimuthal observable to capture the linear polarization.

Introduction.—Pinning down the CP nature of the Higgs-top interaction ($ht\bar{t}$) is an important program being pursued at the Large Hadron Collider (LHC) [1–5]. Any deviation from a Standard-Model-like Higgs-top coupling could indicate new physics as well as provide a potential source for the CP violation as required by the baryogenesis [6]. Unlike CP-violating Higgs interactions with vector bosons, which arise from dimension-six operators, CP-violating effects in the $ht\bar{t}$ coupling could occur via a dimension-four operator,

$$\mathcal{L} \supset -\frac{y_t}{\sqrt{2}} h t (\kappa + i \bar{\kappa} \gamma_5) t,$$  

and can be potentially larger. In Eq. (1), $y_t = \sqrt{2} m_t/v$ is the Yukawa coupling of Higgs and top quark in the Standard Model (SM), and $(\kappa, \bar{\kappa})$ parametrize the CP-even and CP-odd $ht\bar{t}$ interactions, respectively, which can be reparametrized as $(\kappa, \bar{\kappa}) = \kappa_i (\cos \alpha, \sin \alpha)$, with $\alpha$ being the CP phase. The SM corresponds to $(\kappa, \bar{\kappa}) = (1, 0)$ or $(\kappa_i, \alpha) = (1, 0)$.

Numerous approaches have been proposed for determining the CP phase, either directly via associated Higgs and top production [7–20], or indirectly via Higgs or top induced loop effects [21–29]. The sensitivity to $\alpha$ can be enhanced by using observables that are odd under CP transformation [11, 15]. Machine learning techniques have also been considered [17, 20, 30–32] to optimize the sensitivity. The current experimental bounds from direct measurements for various Higgs detection channels are $|\alpha| \leq 35^\circ$ [2], $48^\circ$ [4], and $63^\circ$ [5] at 68% C.L., and $|\alpha| \leq 43^\circ$ at 95% C.L. [5]. These need to be further constrained by more complementary observables, at the upcoming High-Luminosity LHC (HL-LHC) [33] and a possible future pp collider at 100 TeV (FCC-hh) [34].

In this Letter, we propose a new CP-odd observable, for probing the $ht\bar{t}$ interaction, which originates from a linearly polarized gluon in the associated production of a Higgs boson and gluon jet ($hg$). The essential observation is that a singly polarized gluon can be produced from the hard scattering of unpolarized partons. After its production, the gluon fragments into a jet with some linear polarization that breaks the rotational invariance around the jet direction and orients the jet constituents according to the hippopedal distribution,

$$\text{const.} + \xi_1 \cos 2\phi + \xi_2 \sin 2\phi.$$  

Here, as will be defined below, $\xi_1$ and $\xi_2$ parameterize the two degrees of freedom of the linear polarization and depend on both the kinematics of the hard process and the $ht\bar{t}$ couplings, $\kappa$ and $\bar{\kappa}$. The azimuthal angle $\phi$ is defined in the $x\hat{y}$ plane of the coordinate system,

$$\hat{z} = \frac{k}{|k|}, \quad \hat{y} = \frac{\hat{z}_{\text{lab}} \times \hat{z}}{|\hat{z}_{\text{lab}} \times \hat{z}|}, \quad \hat{x} = \hat{y} \times \hat{z},$$  

shown in Fig. 1 (right), where $\hat{z}_{\text{lab}}$ is the beam direction, and $k$ is three-momentum of the gluon jet in the partonic center-of-mass (c.m.) frame.

![Diagram](arXiv:2111.08452v2 [hep-ph] 14 Aug 2024)

FIG. 1. Left: a representative diagram for $gg \to hg$ via a box top loop; the others omitted are $s-, t-$, and $u$-channel diagrams involving triangular top loops and tri-gluon vertices. Right: the gluon $\hat{x}\hat{y}\hat{z}$ frame defined in Eq. (3).

The azimuthally anisotropic jet image in Eq. (2) can be measured as a new jet substructure observable and provide sensitivity to the CP phase of the $ht\bar{t}$ interaction. In particular, we will show that $\xi_2$ and the associated sin $2\phi$ structure are CP-odd. They are more sensitive to a small CP phase $\alpha$ than $\xi_1$, including the sign of $\alpha$. Contrary to building upon a neutral state of charged particles and antiparticles [35], this CP-odd observable is constructed purely out of the kinematic information in the gluon jet. Such CP sensitivity in $hg$ production would not be possible without exploring the gluon jet substructure, which...
has not been considered previously in the literature. We also note that associated Higgs-top production and indirect measurements via $hVV$ or $VV$ production also depend on the $hVV$ couplings and require assumptions on the latter, whereas $hg$ production only depends on the $htt$ coupling.

**Linearly polarized gluon in $hg$ production.**—The polarization state of the produced gluon is described by a density matrix, represented in the helicity basis as

$$\rho_{\lambda\lambda'} = \frac{1}{2} (1 + \xi \cdot \sigma)_{\lambda\lambda'} = \frac{1}{2} \left( 1 + \xi_1 \xi_2 \xi_3 \right),$$

with three polarization degrees of freedom, $\xi \equiv (\xi_1, \xi_2, \xi_3)$. The diagonal element $\xi_3 = \rho_{++} - \rho_{--}$ describes the net gluon helicity, whereas the off-diagonal $\xi_1 = 2 \text{Re} \rho_{+-}$ and $\xi_2 = -2 \text{Im} \rho_{+-}$ are associated with interference between the gluon $+$ and $-$ helicity states. They are better understood in terms of the linear polarization state $|\phi\rangle$ in the $\hat{x}$-$\hat{y}$-$\hat{z}$ frame, related to the helicity eigenstates $|\pm\rangle$ by $|\phi\rangle = [e^{i\phi}|+\rangle - e^{-i\phi}|-\rangle]/\sqrt{2},$

$$\xi_1 = \langle \pi/2 | \rho | \pi/2 \rangle - \langle 0 | \rho | 0 \rangle = \rho_{yy} - \rho_{xx},$$

$$\xi_2 = \langle 3\pi/4 | \rho | 3\pi/4 \rangle - \langle \pi/4 | \rho | \pi/4 \rangle.$$

Thus, $\xi_1$ and $\xi_2$ are differences between linear polarization degrees along two orthogonal directions. It is readily seen that under $CP$ transformation, $(\xi_1, \xi_2) \to (\xi_1, -\xi_2)$ so they are $CP$-even and $CP$-odd, respectively. The ambiguity in defining $\xi_{ab}$ in Eq. (3) at a $pp$ collider merely implies the change $(\hat{x}, \hat{y}) \to (-\hat{x}, -\hat{y})$, which does not affect linear polarization states, contrary to transverse spins of fermions [36].

The gluon produced in the $hg$ process has a significant linear polarization. At leading order (LO), both $gg$ fusion and $gq$ annihilation contribute via a top loop, as exemplified in Fig. 1 (left) for the $gg$ channel. Even though the $gq$ channel can also produce a substantially polarized gluon, its contribution to the total cross section is much smaller and will be neglected. Parametrizing the helicity amplitudes $g(\lambda_1) g(\lambda_2) \to h g(\lambda_3)$ in the partonic c.m. frame in terms of the gluon’s transverse momentum $p_T$, rapidity $y_g$, and azimuthal angle $\phi_g$, we have

$$\mathcal{M}_{\lambda_1,\lambda_2,\lambda_3}(p_T, y_g, \phi_g) = f^{abc} e^{i(\lambda_1 - \lambda_2)\phi_g} \times \kappa \mathcal{A}_{\lambda_1,\lambda_2,\lambda_3}(p_T, y_g) + i \kappa \mathcal{\tilde{A}}_{\lambda_1,\lambda_2,\lambda_3}(p_T, y_g),$$

with $f^{abc}$ the color factor, and $\lambda_i$ the gluon helicities. The $p_T$ and $y_g$ sufficiently determine the Higgs energy, $E_h^2 = m_H^2 + p_T^2 \cosh^2 y_g$, and the partonic c.m. energy $\sqrt{s} = p_T \cosh y_g + E_h$, with $m_H$ being the Higgs mass. $\mathcal{A}$ and $\mathcal{\tilde{A}}$ are the $CP$-even and $CP$-odd helicity amplitudes, respectively, constrained by their $CP$ properties as

$$\mathcal{A}, \mathcal{\tilde{A}}_{-\lambda_1, -\lambda_2, -\lambda_3}(p_T, y_g) = (-\mathcal{A}, +\mathcal{\tilde{A}})_{\lambda_1, \lambda_2, \lambda_3}(p_T, y_g).$$

The gluon density matrix is determined through

$$\frac{1}{4 N_c^2 g} \mathcal{M}_{\lambda_1,\lambda_2,\lambda_3} \mathcal{M}_{\lambda_1,\lambda_2,\lambda_3}^* \equiv \rho_{\lambda\lambda'}(\xi) |\mathcal{M}|^2,$$

where the convention of summing over repeated indices is taken, and $|\mathcal{M}|^2$ is the unpolarized squared amplitude, averaged/summed over the spins and colors, with $N_c = 8$. Due to their $CP$ properties in Eq. (7), $\mathcal{A}$ and $\mathcal{\tilde{A}}$ individually only contribute to $\xi_1$, while it is their interference that contributes to $\xi_2$. In terms of the $CP$ phase $\alpha$, $\xi_1$ and $\xi_2$ can be expressed as

$$\xi_1 = \frac{\omega + \beta_1 \cos 2\alpha}{1 + \Delta \cos 2\alpha}, \quad \xi_2 = \frac{\beta_2 \sin 2\alpha}{1 + \Delta \cos 2\alpha},$$

where we have defined the polarization parameters

$$\Delta = \frac{|\mathcal{A}|^2 - |\mathcal{\tilde{A}}|^2}{|\mathcal{A}|^2 + |\mathcal{\tilde{A}}|^2}, \quad \omega = \frac{2(\mathcal{A}_+ * \mathcal{\tilde{A}}_+ - \mathcal{A}_- * \mathcal{\tilde{A}}_-)}{|\mathcal{A}|^2 + |\mathcal{\tilde{A}}|^2},$$

$$\beta_1 = \frac{2(\mathcal{A}_+ * \mathcal{\tilde{A}}_- - \mathcal{A}_- * \mathcal{\tilde{A}}_+)}{|\mathcal{A}|^2 + |\mathcal{\tilde{A}}|^2}, \quad \beta_2 = \frac{4 \text{Re}(\mathcal{A}_+ * \mathcal{\tilde{A}}_-)}{|\mathcal{A}|^2 + |\mathcal{\tilde{A}}|^2},$$

with the notations

$$A_+ * B_- \equiv A_{\lambda_1,\lambda_2,\lambda_3} B_{\lambda_1,\lambda_2,-\lambda_3}, \quad |\mathcal{A}|^2 \equiv A_{\lambda_1,\lambda_2,\lambda_3} A_{\lambda_1,\lambda_2,\lambda_3}^*.$$

Parametrizing $\xi_{1,2}$ as in Eq. (9) clearly shows that the polarization only depends on the $CP$ phase $\alpha$, but not on the coupling strength $\kappa_t$, which only controls the event rate. The helicity polarization $\xi_3$ is also nonzero as $\sqrt{s} > 2m_t$, but its value is generally small compared to $\xi_1$ and $\xi_2$, and will not be discussed in this work.

The parameters ($\Delta, \omega, \beta_1, \beta_2$) are functions of $p_T$ and $y_g$, as shown in Fig. 2(a) for some benchmark phase-space points. While $\Delta$ describing the relative difference between the $CP$-even and $CP$-odd amplitude squares stays relatively flat around $-0.4$ when $p_T < 10$ TeV, the parameters $\omega, \beta_1, \text{and} \beta_2$, which control the sizes of the polarizations $\xi_1$ and $\xi_2$, vary sizably with $p_T$. Based on their $p_T$ dependence, we divide the phase space into three kinematic regions and discuss them in turn.

1. **Low-$p_T$ region:** $p_T \lesssim 100$ GeV. Both $|\omega|$ and $\beta_1$ have large values, whereas $\beta_2 \approx 0$. The linear polarization is thus dominated by $\xi_1$, with $\xi_2 \approx 0$. The dominance of $\omega$ over $\beta_1$ further implies that $\xi_1$ does not depend sensitively on $\alpha$. Being well below the $\sqrt{s} = 2m_t$ threshold, this region can be well approximated by the infinite-top-mass effective field theory (EFT) [37, 38]. In Fig. 2(b), the SM predictions for $\xi_1$ are shown for both the full one-loop calculation and the EFT approximation, where one can see that $\xi_1$ generally has a large negative value, which means that the produced gluon is dominantly polarized along the $\hat{x}$ direction in the production plane, cf. Eq. (5). Furthermore, it is not dramatically dependent on the gluon rapidity $y_g$.
Since the low-$p_T$ region contains most of the $hg$ events, it is suitable for testing the linear polarization phenomenon. We expect a significant $\cos 2\phi$ jet anisotropy due to the dominant $\xi_1$. The insensitivity to $\alpha$ also enables this region to serve as a calibration region for experimentally measuring the linear polarization, which is important to ensure its viability and to understand the systematic uncertainties of the measurement since such phenomenon has not been observed before.

2. Transition region: 100 GeV $\lesssim p_T \lesssim$ 300 GeV.

The polarized gluon jet function. Around the same time as QCD was developed, it was noted that linearly polarized gluons with nonzero $\xi_1$ can be produced in hard collision processes [39–55], and some non-perturbative arguments were used in favor of oblate gluon jets characterized by a $\cos 2\phi$ distribution. In the presence of a $CP$-violating interaction as considered in this work, a nonzero $\xi_2$ polarization is also produced leading to an additional $\sin 2\phi$ structure, which serves as a handle to probe the $CP$ structure.

Here, we introduce the polarized gluon jet in terms of the modern factorization formalism, for the first time. The polarized gluon turns into a jet that imprints its polarization information in the azimuthal distribution of its constituents, which can be projected out by weighting each event by some azimuth-sensitive observable. The azimuthally weighted cross section $\sigma_w$ of the inclusive $hg$ production at a $pp$ collider can be factorized into a hard scattering cross section, as given in Eq. (8), multiplied by a polarized gluon jet function, in much the same way as the factorization for an unpolarized jet function [56–58] or fragmentation function [59, 60]. It reads as

$$\frac{d\sigma_w}{dy_g \, dp_T^2 \, dm_j^2 \, d\phi} = \frac{d\sigma}{dy_g \, dp_T^2} \frac{dJ(\xi(p_T, y_g), m_j^2, \phi)}{d\phi},$$

up to corrections of powers of $m_j/p_T$ and the jet size $R$. Here, $d\sigma/dy_g \, dp_T^2$ = $\mathcal{L}(s, \hat{s}) |M|^2/16\pi E_h \sqrt{s}$ is the differential cross section for the on-shell gluon production,
where \( L(s, \bar{s}) = \int \frac{d^4 x}{(2\pi)^4} \delta^4(x) f_{q/p}(x, \mu_F) f_{\bar{q}/p}(\bar{s}/x, \mu_F) \)
\
\text{is the gluon-gluon parton luminosity, with the factorization scale chosen at } \mu_F = p_T \text{ in the parton distribution function (PDF) } f_{q/p}(x, \mu_F) \text{ of the proton, and we have used the LO kinematics to integrate over the Higgs phase space.}

In the partonic c.m. frame, the gluon momentum \( k \) defines the jet mass \( m_J^2 = k^2 \) and direction \( \hat{z} \) as in Eq. (3). By defining two lightlike vectors \( n^\mu = (1, -\hat{z})/\sqrt{2} \) and \( \bar{n}^\mu = (1, \hat{z})/\sqrt{2} \), we can approximate the gluon momentum in the hard part to be on shell by only retaining the large component, \( p_T^g = (k \cdot n)\bar{n}^\mu \), which then defines the rapidity \( y_g \) and \( p_T = k \cdot n/(\sqrt{2}\cosh y_g) \). To the leading power of \( m_J/p_T \), the on-shell gluon carries the polarization \( \xi \) and fragments into a jet, described by the polarized jet function \( dJ(\xi, m_J^2, \phi)/d\phi \),

\[
\frac{dJ}{d\phi} = \frac{1}{2\pi N_c(\bar{n} \cdot k)^2} \sum_X \int \frac{d^4 x}{(2\pi)^4} e^{ik \cdot x} \left[ \rho_{\lambda\lambda}(\xi) O(\phi, X) \right] \\
\times \epsilon_{\lambda\nu}(p) \langle 0| W_{a\nu}(\infty, x; n) \sigma G_{\mu}^{\nu*}(x)|X \rangle \\
\times \epsilon_{\lambda\mu}(p) \langle X| W_{ab}(\infty, 0; n) \sigma G_{b}^{\mu}(0)|0 \rangle,
\]

where \( X \) denotes the state of the particles within the jet, in accordance with the jet algorithm [57, 61], whose momenta are dominantly along \( \bar{n} \). \( G_{\mu}^{\nu*} \) is the gluon field strength tensor, and \( W_{ab}(\infty, x; n) \) is the Wilson line in the adjoint representation from \( x \) to \( \infty \) along \( n \), with the color indices \( a, b, \) and \( c \) summed over. In Eq. (12), the gluon polarization vectors \( \epsilon_{\lambda\nu}(p) \) with helicity \( \lambda = \pm 1 \), which are then averaged with the density matrix \( \rho_{\lambda\lambda}(\xi) \). The resultant azimuthal distribution is extracted by inserting the observable

\[
O(\phi, X) = \frac{1}{\sum_{I \in X} p_{i,T} \delta(\phi - \phi_I)},
\]

where \( p_{i,T} \) and \( \phi_I \) are, respectively, the transverse momentum and azimuthal angle of the jet constituent \( i \) with respect to the \( \hat{x} \)-\( \hat{y} \) plane defined in Eq. (3). The \( \phi \) distribution is a new jet substructure observable introduced by the linear polarization. The dependence on \( \xi \) would vanish due to parity invariance of \( O(\phi, X) \).

As a result of the \( p_{i,T} \) weight, the observable \( O(\phi, X) \) is infrared (IR) safe, and hence the polarized gluon jet function is insensitive to hardonization effects and becomes perturbatively calculable, with a predictable \( \phi \) dependence. However, it was noted long before [43, 50] that the gluon polarization information will be greatly washed out by the cancellation between the \( g \to gg \) and \( g \to g\bar{q} \) channels, which was also found recently in a similar situation [62–64]. It is possible to mitigate these effects by using jet flavor tagging techniques [65–76]. For example, one may recluster the identified gluon jet into two subjets, and only keep those gluon jets with their two subjets tagged as quarks. At \( O(\alpha_s) \), requiring a tagged quark in the gluon jet leaves \( g \to q\bar{q} \) as the only diagram, giving the polarized gluon jet function,

\[
\frac{dJ}{d\phi} = \frac{a_s T_F}{6\pi^2 m_J^2} \left[ 1 + \frac{1}{2} \left( \xi_1 \cos 2\phi + \xi_2 \sin 2\phi \right) \right],
\]

where the jet algorithm dependence does not come in at this order to the leading power of \( m_J \). Eq. (14) needs to be multiplied by the tagging efficiency when used in Eq. (11). Although flavor tagging reduces the statistics significantly, it enhances the gluon spin analyzing power from \( O(1\%) \) to about 50% [50] and will improve the statistical precision.

Before closing this section, we note the difference of the gluon polarization from a quark. While a transversely polarized light (massless) quark can also be produced from hard scattering processes, its transverse spin cannot be conveyed via the perturbative quark jet function due to the chiral symmetry of a massless quark. It is hence related to chiral symmetry breaking and must require the presence of some non-perturbative functions [60, 77, 78].

**Phenomenology.**—The gluon jet azimuthal anisotropy in Eq. (14) can be experimentally measured by simply constructing the asymmetry observables [79]

\[
A_i = \frac{\int_0^{2\pi} d\phi \langle d\sigma_w/d\phi \rangle \cdot \text{sgn}[F_i(\phi)]}{\int_0^{2\pi} d\phi \langle d\sigma_w/d\phi \rangle} = \frac{\xi_i}{\pi},
\]

where \( i \in \{1, 2\} \), \( F_i(\phi) = \cos 2\phi \) and \( F_2(\phi) = \sin 2\phi \). The uncertainties of the asymmetries \( A_{1,2} \) are dominated by statistical ones, given by \( 1/\sqrt{N} \) with \( N \) being the number of the observed events. Now we provide a simple demonstration of the constraining power of the gluon linear polarization on the CP phase, by confining ourselves to the transition region for both the HL-LHC at 14 TeV and FCC-hh at 100 TeV, with integrated luminosities 3 ab\(^{-1}\) and 20 ab\(^{-1}\), respectively.

The \( hg \) cross section in the transition region is estimated for the Lagrangian [Eq. (1)] using CT18NNLO PDFs [80] with MGS_aMC@NLO 2.6.7 [81] by first generating the \( hg \) events with \( p_T \in [100, 300] \) GeV and \( |y_g| \leq 2.5 \) in the lab frame, and then boosting to the partonic c.m. frame with a further cut \( |y_g| \leq 0.8 \), which gives \( \sigma_T^g(0.57 \cos^2 \alpha + 1.3 \sin^2 \alpha) \) pb for the HL-LHC and \( \sigma_T^g(13.7 \cos^2 \alpha + 30.7 \sin^2 \alpha) \) pb for the FCC-hh. While both \( \kappa_4 \) and \( \alpha \) affect the total production rate and can be constrained by the measurement of the latter, only \( \alpha \) determines the polarization. In the following, we take \( \kappa_4 = 1 \) and consider the constraint on \( \alpha \) from the polarization data.

We are interested in final states where the (fat) gluon jet is composed of a pair of quark subjets. While it is possible to also discriminate light quark subjets from gluon subjets, here we only provide a conservative estimate by restricting to the bottom (b) and charm (c) quark tagging as used in experiments [82–92]. We estimate the branching fraction \( f_{gqq} (f_{gcc}) \) of \( g \to b\bar{b} (g \to c\bar{c}) \) through parton
shower simulation using Pythia 8.307 [93], which gives $f_{h_{bb}} = 0.013$ and $f_{h_{cc}} = 0.019$ in the selected kinematic region. Following Refs. [86, 87], we take $b$-tagging efficiency $\epsilon_{b} = 0.7$ and $c$-tagging efficiency $\epsilon_{c} = 0.3$. We consider the diphoton decay channel of the SM Higgs boson and assume a Higgs tagging efficiency $\epsilon_{t} = 0.002$. This then gives about $(51 \cos^{2} \alpha + 115 \sin^{2} \alpha)$ reconstructed events at the HL-LHC and $(8100 \cos^{2} \alpha + 18200 \sin^{2} \alpha)$ events at the FCC-hh.

In Fig. 3, we display the predicted average values of $\xi_{1,2}$ in the transition region at the FCC-hh as functions of the CP phase $\alpha$. $(\xi_{1,2})$ are the average values of $\xi_{1,2}$ in the specified kinematic region. Their statistical uncertainties are indicated by the red and blue bands, respectively, around the SM prediction (with $\alpha = 0$). The green-hatched region is the $\alpha$ range allowed by the $\xi_{2}$ measurement.

**Summary.**—A precise understanding of the CP property of the Higgs boson is important both to test the SM and to probe new physics. In this Letter, we proposed a novel way of probing the CP structure of the Higgs-top interaction, by measuring the azimuthal anisotropy substructure of the gluon jet produced in association with a Higgs boson, which originates from the linear polarization of the final-state gluon. We have introduced a factorization formalism and defined a perturbative polarized gluon jet function with insertion of an IR-safe azimuthal observable. Experimental measurement of the linearly polarized gluon jet will be an important test of the SM and can also serve as a new tool to search for new physics.

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