QCD Corrections to Jet Correlations in Weak Boson Fusion

Terrance Figy and Dieter Zeppenfeld

*Department of Physics, University of Wisconsin, Madison, WI*

**Abstract**

Higgs boson production via weak boson fusion is sensitive to the tensor structure of the $HVV$ ($V = W, Z$) couplings, which distinguishes loop induced vertices from SM expectations. At the CERN Large Hadron Collider this information shows up most clearly in the azimuthal angle correlations of the two forward and backward quark jets which are typical for weak boson fusion. We calculate the next-to-leading order QCD corrections to this process, in the presence of anomalous $HVV$ couplings. Gluon emission does not significantly change the azimuthal jet correlations.
**Introduction.** The production of Higgs bosons in the weak boson fusion (WBF) process will provide a direct and highly sensitive probe of $HWW$ and $HZZ$ couplings at the CERN Large Hadron Collider (LHC) \cite{1,2,3,4,5}. The determination both of the strength and of the tensor structure of these couplings is crucial for the identification of the produced boson as a remnant of the spontaneous symmetry breaking process which is responsible for $W$ and $Z$ mass generation.

Within spontaneously broken, renormalizable gauge theories like the standard model (SM), this coupling originates from the kinetic energy term, $(D_\mu \Phi)^\dagger (D^\mu \Phi)$, of a scalar Higgs field, $\Phi$, whose neutral component obtains a vacuum expectation value (vev), $\Phi^0 \rightarrow (v + H)/\sqrt{2}$. This replacement then leads to a characteristic coupling in the interaction Lagrangian, of the form $HV_\mu V^\mu$ ($V = W, Z$). The existence of the vev is necessary to produce a trilinear $HVV$ coupling at tree level: with $v = 0$ all couplings to the gauge fields $V$ contain two scalar fields, i.e., only $HHV$ and $HHVV$ couplings would be generated. A trilinear $HVV$ coupling may also be loop-induced, however. The SM $H\gamma\gamma$ and $Hgg$ effective couplings are an example: they are induced by $W$-boson and/or top quark loops. Gauge invariance dictates a different tensor structure of these loop-induced couplings: the corresponding effective Lagrangian contains the square of the field strength, i.e. the lowest order loop-induced terms are of the form $HV_{\mu\nu}V^{\mu\nu}$ or $HV_{\mu\nu}\tilde{V}^{\mu\nu}$, where $\tilde{V}^{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}V_{\rho\sigma}$ denotes the dual field strength of the gauge field.

The task of future Higgs experiments is, then, twofold: (i) to measure the overall strength of the $HVV$ coupling, and (ii) to identify its tensor structure. One would expect a loop-induced coupling to be much smaller than the expected SM $HVV$ coupling strength. However, the measurement of WBF rates alone will not be sufficient to establish $H$ as being related to spontaneous symmetry breaking: to give just two examples, the loop-induced couplings might be substantially enhanced by additional non-SM particles in the loop or by the existence of multiplets of large weak isospin which couple strongly to $H$. Or a particular LHC signature may be strongly enhanced by a much larger $H$ decay branching ratio than in the SM. A confirmation that the $HVV$ coupling has tree level strength is, thus, ambiguous: a clear identification of the Higgs boson also requires the identification of the tensor structure of the $HVV$ vertex.

It was pointed out some time ago that the azimuthal angle correlations of the two quark jets in the weak boson fusion process $qQ \rightarrow qQH$ provide tell-tale signatures for the tensor
structure of the $HVV$ couplings $\phi_{jj}$: the SM expectation is for a flat distribution, while the loop-induced couplings lead to a pronounced dip at azimuthal separations $\phi_{jj}$ of the two tagging jets of 90 degrees for a $HV_{\mu\nu}V^{\mu\nu}$ coupling and at 0 and 180 degrees for the CP violating $HV_{\mu\nu}\bar{V}^{\mu\nu}$ vertex. Observation of the tagging jets is crucial for isolating the WBF process from backgrounds and, therefore, their distributions will be available for all WBF samples. Also, signal to background ratios for WBF processes are expected to be very good within the SM, exceeding the 1:1 level for wide ranges of the Higgs boson mass $m_H$.

The analysis of Ref. [6] was performed at leading order (LO) in QCD. This means that additional gluon emission, which might lead to a de-correlation of the tagging jets, was ignored in the analysis. Subsequently it was argued [7] that such de-correlation effects play an important role in a related process, $gg \rightarrow Hgg$, when the two tagging jets are widely separated in rapidity, which is a typical requirement for WBF studies. In this Letter we analyze this question, by calculating the tagging jet distributions in next-to-leading order (NLO) QCD, for the production of a scalar $H$ via WBF with an arbitrary tensor structure of the $HVV$ vertex. If de-correlation is important, it should show up in the form of large radiative corrections at NLO. We use the term “Higgs boson” as a generic name for the produced scalar in the following.

The NLO calculation. Our calculation is an extension of the NLO QCD corrections for the SM WBF processes $qQ \rightarrow qQH$ (and crossing related ones) [8, 9, 10]. For the total cross section these corrections have been known for over a decade [8]. Recently, we have recalculated them by developing a NLO parton level Monte Carlo program [9], which provides the flexibility to calculate arbitrary distributions at NLO, such as the azimuthal angle correlations that we are interested in here.

The calculation of Ref. [9] uses a SM vertex function, $T^{\mu\nu}(q_1, q_2) = \frac{2m^2_v}{v} g^{\mu\nu}$ for the $HVV$ vertex in Fig. 1. Here we need to generalize this vertex to the most general structure compatible with Lorentz invariance. Taking into account that the quark currents in Fig. 1 and for the corresponding gluon emission processes are conserved, all terms proportional to $q_1^\mu$ or $q_2^\nu$ may be dropped, and the most general $HVV$ vertex may be written as

$$T^{\mu\nu}(q_1, q_2) = a_1(q_1, q_2) g^{\mu\nu} + a_2(q_1, q_2) [q_1 \cdot q_2 g^{\mu\nu} - q_2^{\mu} q_1^{\nu}] + a_3(q_1, q_2) \varepsilon^{\mu\nu\rho\sigma} q_1^\rho q_2^\sigma. \tag{1}$$

Here $q_1$ and $q_2$ are the four-momenta of the two weak bosons, and the $a_i(q_1, q_2)$ are Lorentz-invariant form factors, which might, for example, represent scalar loop integrals in a per-
FIG. 1: Feynman graphs contributing to $\bar{q}Q \rightarrow \bar{q}QH$ at (a) tree level and (b) including virtual corrections to the upper quark line. The momentum labels and Lorentz indices for the internal weak bosons correspond to the vertex function of Eq. (1).

In our discussion we will neglect possible contributions from $H\gamma\gamma$ and $H\gamma Z$ couplings which
can appear in $SU(2) \times U(1)$ invariant formulations \cite{13, 14}. The precise mix of $HWW$, $HZZ$, $HZ\gamma$ and $H\gamma\gamma$ contributions is quite irrelevant for the observable azimuthal angle distributions, as long as we do not consider interference effects between SM and anomalous vertices, and it will not affect our conclusions about the size of NLO corrections. For simplicity we therefore set $a_1 = 0$ for the anomalous coupling case and choose relative contributions from $WW$ and $ZZ$ fusion as in the SM, by taking $g_{5o}^{HWW} = g_{5e}^{HWW} = 1$, $g_{5e}^{HZZ} = g_{5o}^{HZZ} = 1/\cos^2 \theta_W$, and by using either $\Lambda_{5e} \simeq 480$ GeV, $\Lambda_{5o} = \infty$ for the CP even case or $\Lambda_{5o} \simeq 480$ GeV, $\Lambda_{5e} = \infty$ for the CP odd case, which roughly reproduces SM rates for a scalar mass of $m_H = 120$ GeV.

The effective Lagrangian of Eq. (2) produces couplings

$$a_2(q_1, q_2) = -\frac{2}{\Lambda_{5e}} g_{5e}^{HWW}, \quad a_3(q_1, q_2) = \frac{2}{\Lambda_{5o}} g_{5o}^{HWW}$$

for the $HWW$ vertex, and

$$a_2(q_1, q_2) = -\frac{2}{\Lambda_{5e}} g_{5e}^{HZZ}, \quad a_3(q_1, q_2) = \frac{2}{\Lambda_{5o}} g_{5o}^{HZZ}$$

for the $HZZ$ vertex. In general, the $a_i$ are form factors which are expected to be suppressed once the momentum transfer, $\sqrt{-q_i^2}$, carried by the virtual gauge boson reaches the typical mass scale, $M$, of the new physics which is responsible for these anomalous couplings. Below we use the simple ansatz

$$a_i(q_1, q_2) = a_i(0, 0) \frac{M^2}{q_1^2 - M^2} \frac{M^2}{q_2^2 - M^2}$$

for discussing the consequences of such form factor effects.

Results: The typical signature of a weak boson fusion event at the LHC consists of the two quark jets (tagging jets) and the Higgs decay products. The tagging jets tend to be widely separated in rapidity, with one quite forward (typical pseudorapidity of 3 to 4) and the second one backward, but frequently still located in the central detector (pseudorapidity below 2.5). Various Higgs decay modes have been considered in the literature for WBF, $H \rightarrow WW$ \cite{1}, $H \rightarrow \tau\tau$ \cite{2}, and $H \rightarrow \gamma\gamma$ \cite{3} being the most promising ones. While optimized event selection varies, in particular for the decay products, the cuts on the tagging jets are fairly similar in all analyses. Since here we are interested in the QCD features of WBF events, which do not depend on the Higgs decay mode, we perform our NLO analysis without simulating Higgs decays, and we only impose typical WBF cuts on the tagging jets.
FIG. 2: Normalized transverse momentum distribution of the hardest jet for the SM Higgs boson (solid red line) and a scalar $H$ of mass $m_H = 120$ GeV with CP even anomalous coupling $a_2(q_1, q_2)$. The dash-dotted curves correspond to different form factor scales $M = 100, 200, 400$ GeV in Eq. (5) and $a_2 = \text{const.}$ (blue curves) at NLO. LO curves are shown by the dashed lines and differ very little from the NLO results.

In order to reconstruct jets from the final-state partons, the $k_T$-algorithm [15] as described in Ref. [16] is used, with resolution parameter $D = 0.8$. In a given event, the tagging jets are then defined as the two jets with the highest transverse momentum, $p_T j$, with

$$p_T j \geq 20 \text{ GeV}, \quad |y_j| \leq 4.5.$$  \hspace{1cm} (6)

Here $y_j$ denotes the rapidity of the (massive) jet momentum which is reconstructed as the four-vector sum of massless partons of pseudorapidity $|\eta| < 5$. Backgrounds to weak-boson fusion are significantly suppressed by requiring a large rapidity separation of the two tagging jets. This motivates the final cut

$$\Delta y_{jj} = |y_{j_1} - y_{j_2}| > 4, \quad y_{j_1} \cdot y_{j_2} < 0,$$  \hspace{1cm} (7)

which includes the requirement that the two tagging jets reside in opposite detector hemispheres.
The structure of the $HVV$ coupling affects the production dynamics of $H$ and we can expect significant deviations in jet observables if, instead of the SM, anomalous couplings describe the vertex of Eq. (1). One example is shown in Fig. 2 where transverse momentum distributions, $d\sigma/dp_{Tj}(\text{max})$, are compared between the SM (solid line) and the CP even coupling $a_2(q_1, q_2)$, with different form factor scales $M$ in Eq. (5). Here, $p_{Tj}(\text{max})$ is the maximum $p_T$ of the two tagging jets. Only the shape of the distribution is considered, since the rate can always be adjusted by multiplying the anomalous couplings by a constant factor. Also, we should note that a CP odd coupling leads to very similar curves for a given form factor scale. In all cases we show the LO expectations (dashed lines) together with the NLO results: QCD corrections are of order 10%, typically, and well under control.

One finds that anomalous $HVV$ couplings generally lead to harder $p_T$ spectra of the two tagging jets. Since the anomalous Lagrangian in Eq. (2) couples the Higgs boson to weak boson field strengths, transverse polarizations of the incident $V V$ pairs dominate the anomalous case, while longitudinal $V V$ fusion is responsible for SM Higgs production. A telltale sign of transverse vector boson fusion is the more central and, hence, higher $p_T$ production of the tagging jets. This effect is enhanced by the momentum factors in the $HVV$ anomalous vertices.

While the changed transverse momentum distributions in Fig. 2 could be used to rule out the SM, the reverse is not readily possible: a jet transverse momentum distribution compatible with SM expectations might be faked by anomalous couplings and a judiciously chosen form factor behavior of the coefficient functions $a_2$ or $a_3$ in Eq. (5). The different scale choices in Fig. 2 demonstrate this effect: a low form factor scale of $M = 100$ GeV or slightly lower would be difficult to distinguish from the SM expectation and one can certainly find a functional form of the form factors which reproduces the SM within experimental errors.

A much better observable for distinguishing the different tensor structures of the $HVV$ vertex is the azimuthal angle correlation of the two tagging jets, $d\sigma/d\phi_{jj}$ [6]. Here $\phi_{jj}$ is the azimuthal angle between the two tagging jets. The corresponding distributions are shown in Fig. 3 for the SM (solid line) and for the same choices of form factors as before. The dip at $\phi_{jj} = 90$ degrees for the CP even coupling and the suppression at 0 and 180 degrees for the CP odd coupling are clean signatures which only depend on the tensor structure of the couplings and not on the precise dynamics which is responsible for the form factors. The remaining form factor dependence is very small and can be explained by kinematic effects.
related to the higher average jet transverse momentum for big form factor scales, $M$: at small $\phi_{jj}$ two high $p_T$ jets recoil against the $H$ scalar, resulting in an increased invariant mass of the event compared to the situation with two back-to-back jets. This leads to a more asymmetric $\phi_{jj}$ distribution for high form factor scales.

The pronounced dip at 90 degrees, which is characteristic of the CP even coupling, is also found in $Hjj$ production via gluon fusion $^{17}$, at LO. This is not surprising because, in the large top mass limit, the $Hgg$ vertex can be described by an effective Lagrangian proportional to $HG_a^{\mu\nu}G^{a\mu\nu}$, which exhibits the same field strength squared behavior and hence the same tensor structure as the CP even $HVV$ coupling in Eqs. $^{12}$. Since the two tagging jets are far apart from each other, separated by a large rapidity gap of 4 units of rapidity or more, this LO behavior may be significantly reduced by gluon radiation when higher order QCD corrections are taken into account. Such de-correlation effects have been studied for dijet events at the Tevatron $^{18}$. For $Hjj$ production via gluon fusion, Odagiri $^{7}$ has argued that the dip structure is largely washed out by additional gluon emission between the two tagging jets.
FIG. 4: Higgs mass dependence of the azimuthal angle separation $\phi_{jj}$ of the two tagging jets. In (a) the normalized azimuthal angle distributions are shown at LO (dashed lines) and NLO (solid lines) for Higgs masses of $m_H = 120, 200, 500$ GeV and a constant CP even anomalous coupling $a_2$. Corresponding K-factors are shown in (b).

Our NLO calculations show that such de-correlation effects are irrelevant for weak boson fusion, where $t$-channel color singlet exchange severely suppresses gluon radiation in the central region. The LO and the NLO curves in Fig. 3 are virtually indistinguishable. In order to better exhibit the size of NLO QCD effects for the WBF case, we show, in Fig. 4(a) the azimuthal angle correlations for a pure CP even anomalous coupling for three different Higgs masses, $m_H = 120, 200$ and 500 GeV. Only small changes are visible when going from LO (dashed lines) to NLO (solid lines). The differences between LO and NLO are smaller than kinematical effects that can be induced by cuts on the Higgs decay products or by variations of the Higgs boson mass.

The small to modest size of the QCD corrections is quantified in Fig. 4(b) where the K-factor for the distribution is shown, which is defined as

$$K(\phi_{jj}) = \frac{d\sigma^{NLO}/d\phi_{jj}}{d\sigma^{LO}/d\phi_{jj}}.$$  \hspace{1cm} (8)

The $K$-factor is below $\approx 1.4$ even in the dip region, where the cross section is severely suppressed. Virtually identical results hold for the CP-odd case. Clearly, the characteristic azimuthal angle distributions of the jets in WBF are not affected in any significant way by NLO QCD corrections.
Conclusions: We have performed a first calculation of the NLO QCD corrections to Higgs boson production via WBF in the presence of arbitrary anomalous $HVV$ ($V = W, Z$) couplings. Anomalous couplings lead to characteristic changes in the azimuthal angle correlation of the two tagging jets in weak boson fusion events at the LHC, which provides for a very sensitive test of the tensor structure of the $HVV$ couplings of the Higgs boson or of any other scalar with sufficiently large production cross section in WBF [6]. We have shown by explicit calculation that these azimuthal correlations are not washed out by gluon emission, at NLO QCD, even though the tagging jets are widely separated in rapidity. This behavior can be understood as a consequence of $t$-channel color singlet exchange in WBF which severely suppresses the central gluon radiation which might cause tagging jet de-correlation.

Acknowledgments

This research was supported in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation and in part by the U.S. Department of Energy under Contract No. DE-FG02-95ER40896.

[1] D. Rainwater and D. Zeppenfeld, Phys. Rev. D 60, 113004 (1999) [Erratum-ibid. D 61, 099901 (2000)] arXiv:hep-ph/9906218; N. Kauer, T. Plehn, D. Rainwater and D. Zeppenfeld, Phys. Lett. B 503, 113 (2001) arXiv:hep-ph/0012351; C. M. Buttar, R. S. Harper and K. Jakobs, ATL-PHYS-2002-033; K. Cranmer et al., ATL-PHYS-2003-002 and ATL-PHYS-2003-007; S. Asai et al., ATL-PHYS-2003-005.

[2] D. Rainwater, D. Zeppenfeld and K. Hagiwara, Phys. Rev. D 59, 014037 (1999) [arXiv:hep-ph/9808468]; T. Plehn, D. Rainwater and D. Zeppenfeld, Phys. Rev. D61, 093005 (2000) [arXiv:hep-ph/9911385]; S. Asai et al., ATL-PHYS-2003-005.

[3] D. Rainwater and D. Zeppenfeld, JHEP 9712, 005 (1997) [arXiv:hep-ph/9712271]; K. Cranmer, B. Mellado, W. Quayle and S. L. Wu, arXiv:hep-ph/0401088

[4] S. Asai et al., arXiv:hep-ph/0402254

[5] D. Zeppenfeld, R. Kinnunen, A. Nikitenko and E. Richter-Was, Phys. Rev. D62, 013009 (2000) arXiv:hep-ph/0002036; D. Zeppenfeld, in Proc. of the APS/DPF/DPB Summer
Study on the Future of Particle Physics (Snowmass 2001) ed. N. Graf, eConf C010630, P123 (2001) \texttt{arXiv:hep-ph/0203123}; A. Belyaev and L. Reina, JHEP 0208, 041 (2002) \texttt{arXiv:hep-ph/0205270}.

[6] T. Plehn, D. Rainwater and D. Zeppenfeld, Phys. Rev. Lett. 88, 051801 (2002) \texttt{arXiv:hep-ph/0105325}.

[7] K. Odagiri, JHEP 0303, 009 (2003) \texttt{arXiv:hep-ph/0212215}.

[8] T. Han, G. Valencia and S. Willenbrock, Phys. Rev. Lett. 69, 3274 (1992) \texttt{arXiv:hep-ph/9206246}.

[9] T. Figy, C. Oleari and D. Zeppenfeld, Phys. Rev. D 68, 073005 (2003) \texttt{arXiv:hep-ph/0306109}.

[10] E. L. Berger and J. Campbell, \texttt{arXiv:hep-ph/0403194}.

[11] K. Hagiwara and D. Zeppenfeld, Nucl. Phys. B 274, 1 (1986); K. Hagiwara and D. Zeppenfeld, Nucl. Phys. B 313, 560 (1989).

[12] J. Pumplin et al., JHEP 0207, 012 (2002) \texttt{arXiv:hep-ph/0201195}.

[13] W. Buchmüller and D. Wyler, Nucl. Phys. B 268, 621 (1986).

[14] K. Hagiwara, S. Ishihara, R. Szalapski and D. Zeppenfeld, Phys. Rev. D 48, 2182 (1993).

[15] S. Catani, Yu. L. Dokshitzer and B. R. Webber, Phys. Lett. B 285 291 (1992); S. Catani, Yu. L. Dokshitzer, M. H. Seymour and B. R. Webber, Nucl. Phys. B 406 187 (1993); S. D. Ellis and D. E. Soper, Phys. Rev. D48 3160 (1993).

[16] G. C. Blazey et al., \texttt{arXiv:hep-ex/0005012}.

[17] V. Del Duca, W. Kilgore, C. Oleari, C. Schmidt and D. Zeppenfeld, Phys. Rev. Lett. 87, 122001 (2001) \texttt{arXiv:hep-ph/0105129}; Nucl. Phys. B 616, 367 (2001) \texttt{arXiv:hep-ph/0108030}.

[18] A. H. Mueller and H. Navelet, Nucl. Phys. B 282, 727 (1987); V. Del Duca and C. R. Schmidt, Phys. Rev. D 49, 4510 (1994) \texttt{arXiv:hep-ph/9311290}; W. J. Stirling, Nucl. Phys. B 423, 56 (1994) \texttt{arXiv:hep-ph/9401266}; S. Abachi et al. [D0 Collaboration], Phys. Rev. Lett. 77, 595 (1996) \texttt{arXiv:hep-ex/9603010}; L. H. Orr and W. J. Stirling, Phys. Rev. D 56, 5875 (1997) \texttt{arXiv:hep-ph/9706529}; J. Kwiecinski, A. D. Martin, L. Motyka and J. Outhwaite, Phys. Lett. B 514, 355 (2001) \texttt{arXiv:hep-ph/0105039}.

[19] The dimension 5 language is appropriate for, e.g., an isosinglet scalar resonance $H$. For a Higgs doublet $\Phi$ with a vev, the leading operators appear at dimension 6 level \cite{13,14} and the couplings in Eq. 2 are suppressed by an additional factor $g_5^{HVV} \sim v/\Lambda$. 

11