Disentangling Strong Dynamics through Quantum Interferometry Postprint

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Abstract

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Full Text

Preamble

Disentangling Strong Dynamics through Quantum Interferometry

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We present a new probe of strongly coupled electroweak symmetry breaking at the 14 TeV LHC by measuring a phase shift in the event distribution of

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Introduction

One of the most important goals of the Large Hadron Collider (LHC) is to determine the nature of the electroweak symmetry breaking (EWSB) mechanism. The discovery of a 125 GeV Higgs-like object [**?**, **?**] represents a major milestone in this direction. However, to truly understand EWSB, we must learn the origin of the longitudinal components of massive electroweak gauge bosons (V). Broadly speaking, models of EWSB fall into two categories: those where the longitudinal components of the V's are (a) weakly interacting or (b) strongly interacting. The most popular examples in the first category are the Standard Model (SM) with one or more elementary Higgs multiplets [**?**] and its supersymmetric extensions [**?**]. In the second category, one or more strongly interacting sectors appear at the TeV scale and are responsible for EWSB. The Higgs-like boson and the longitudinal V's could arise as pseudo-Nambu-Goldstone bosons (PNGBs) in this scenario [**?**]. A definitive distinction between these two cases at the LHC would be extremely important and would serve as one of the first crucial steps toward a full understanding of the EWSB mechanism.

A universal consequence of a strongly coupled EWSB sector is an enhancement of the longitudinal gauge boson scattering amplitude at high energies beyond the Standard Model expectation [**?**]. To measure such an effect, one would typically look for an enhanced rate for total gauge boson scattering at high invariant mass for the gauge boson pair, and then use the polar angle distribution of the gauge boson decay products to measure their longitudinal polarization fraction [**?**]. In addition to an enhancement of the amplitude's magnitude, one also expects a large phase shift, as in the case of scattering through a resonance. In this paper, we seek strategies to experimentally probe this phase shift induced by strong dynamics. We use the azimuthal angle correlations of the V's decay products and show that the strong phase shift appears as a modification to the interference effect between the longitudinal and transverse V polarizations in the azimuthal angle distributions. Thus, the phase shift becomes an observable that can serve as a complementary probe (in addition to the longitudinal V scattering rate) of EWSB from strong dynamics.

PNGB Scattering and Parameterization of the Phase Shift

Pion scattering provides a realistic example of strongly coupled PNGB scattering. Experimental $\pi\pi$ scattering data [**?**] reveal a large phase shift in the form factor of both $\pi\pi \to \pi\pi$ and $e^+e^- \to \pi\pi$. The phase shift δ in $e^+e^- \to \pi\pi$

scattering versus energy s is shown by the black data points in Fig. 1 [Figure 1: see original paper]. The amplitude undergoes a large phase shift when s is near the mass of the   meson (760 MeV).

The low-energy effective theory of pion scattering is described by a chiral Lagrangian. In this description, the scattering amplitude in the isospin (I) and partial wave (J) channel of $\pi\pi \to \pi\pi$ is given by M_{IJ}   sin(δ{IJ})e^{iδ{IJ}}   c_{IJ} e^{is/f_π^2}, where f_π = 84 MeV is a low-energy constant of chiral perturbation theory (which differs from the physical pion decay constant, f_π^{phys} = 92 MeV that arises away from the chiral limit), and the last equality follows from the low-energy theorem prediction of the scattering amplitude's behavior with s, which also determines the constants c_{IJ}. Naive extrapolation of this form violates unitarity for energies s comparable to the cutoff 4πf_π of the effective theory.

There are several ways to incorporate the effect of vector resonances into the scattering amplitude calculations of the low-energy theory. One approach is to add a broad vector resonance by hand to the theory. However, we adopt a different approach that manifestly maintains the unitarity of the amplitude at very high energies and emphasizes the central role of the phase shift. We multiply the tree-level amplitude (in the J = 1 channel) by a complex form factor F(s). The entire form factor can then be extracted from its phase using analyticity arguments to define an Omnès function [**?**] and assuming no inelastic channels. For a given form factor F(s), applying the subtracted dispersion relation to log(F(s))/s yields:

$$F(s) = P(s)\exp\left[\int_0^\infty ds' \frac{\delta(s')}{\pi}\left(\frac{1}{s'-s-i\epsilon} - \frac{1}{s'}\right)\right].$$

When the phase δ(s') goes beyond 2π (for instance, with multiple resonances and additional branches), the additional 2πn phase factors can be recast into P(s) as a polynomial factor with P(0) = 1.

For simplicity, we use an ansatz for the phase:

$$\delta(s) = \begin{cases} \arctan\left[\frac{s\Gamma}{m(m^2+\Gamma^2-s)}\right], & s < m^2 \\ \arctan\left[\frac{s\Gamma}{m(m^2+\Gamma^2-s)}\right] + \pi, & s \geq m^2, \end{cases}$$

which approaches a constant at high energy according to unitarity. This ansatz fits the phase of the $\pi\pi$ scattering data very well, as shown in Fig. 1. Our parametrization is general and does not rely on the specific form of strong dynamics. From the phase δ(s), we can construct the full form factor:

$$F(s) = P(s)\exp\left[\frac{1}{\pi}\int_0^\infty ds'\,\delta(s')\left(\frac{1}{s'-s-i\epsilon} - \frac{1}{s'}\right)\right].$$

For P(s) = 1, the magnitude |F(s)| behaves as follows: at small s, far from the resonance, the form factor is unity; as we approach the resonance, the magnitude grows large; for large s beyond the resonance, the form factor falls off rapidly with energy.

Longitudinal Weak Boson Scattering

In a strongly interacting EWSB theory, the longitudinal components of weak bosons can be approximately regarded as PNGBs, and their interactions can be described at low energies by a chiral Lagrangian. Thus, naive extrapolation of the scattering amplitude would lead to a similar unitarity violation problem. We expect a similar resolution as in the pion case, where new resonances unitarize the scattering amplitude.

However, there is one key difference between weak boson scattering and pion scattering: we have already discovered a 125 GeV Higgs-like scalar object that couples to weak bosons. The exchange of this scalar partially restores unitarity in the longitudinal V scattering amplitude. If the couplings of this scalar to V's are exactly the Standard Model value, corresponding to a non-composite Higgs boson, then complete unitarity restoration would occur just from exchange of this scalar.

For composite Higgs models, we can use the generalized Adler-Weinberg sum rule (in the limit of vanishing gauge couplings) [?] to relate the Higgs coupling to V's to an integral sum of longitudinal gauge boson scattering cross-sections in various isospin channels:

$$1 - a^2 = \int_0^\infty ds \left(2\sigma_{\rm tot}^{I=0}(s) + 3\sigma_{\rm tot}^{I=1}(s) - 5\sigma_{\rm tot}^{I=2}(s)\right).$$

Here, a parameterizes the ratio of the Higgs boson coupling to V's over its SM value. From the latest fits [?], we see that $0.8 < a < 1.2$ at the 2σ level.

For $a \neq 1$, we need additional contributions from strong dynamics to restore unitarity in longitudinal gauge boson scattering. For simplicity, we assume vector meson dominance for the remaining partial unitarity restoration, as in pion scattering, restricting the new physics contribution to the J = 1 channel.

We parameterize strong dynamics in longitudinal V scattering by introducing a form factor that multiplies the amplitude, just as we did for $\pi\pi$ scattering. We assume the same ansatz for the phase, since it manifestly maintains unitarity. The role of the 125 GeV boson in unitarization can be accounted for without explicitly including Higgs exchange diagrams. Instead, we note that the J = 1, longitudinal V cross-section must be rescaled by a factor of $1 - a^2$ (compared to the case where no Higgs-like boson is present). Thus, we simply rescale our form factor by $1 - a^2$ to account for the presence of the 125 GeV boson.

Assuming a = 0.8, we show the currently excluded region for different form-factor parameters m and Γ in Fig. 2 [Figure 2: see original paper] using the latest LHC searches for W' → WZ resonances [?]. This search strategy only probes the enhancement to the total rate for WZ production, or equivalently it probes |F(s)|, but it does not probe the phase shift directly. In addition, the search uses a narrow WZ invariant mass window to suppress backgrounds and is therefore insensitive to very broad resonances with Γ/m 20%, which are characteristic of strong dynamics.

Observation of Longitudinal V Scattering

There are two methods to observe VV scattering at the LHC. One is the widely used weak boson fusion process pp → VVjj, where the two forward jets suppress large SM backgrounds. The other is the rescattering process pp → VV, which is not considered promising for most studies due to large SM backgrounds. However, this channel is not suppressed by the small effective-V luminosity and has better access to higher energies of the VV system. Since we aim to observe the azimuthal angle correlation arising from a quantum interference term, we do not need to suppress SM backgrounds and will consider the rescattering process in this paper.

Angular Correlation

Let us consider W^+Z production from a $u\bar{d}$ initial state. We modify the SM amplitude of longitudinal W^+Z production in the J = 1 channel by a complex form-factor (1 - a^2)F(s), where F(s) is given by the form in Eq. (4). We study this process at high energies where both the W^+ and Z decay leptonically. The kinematic dependence of the production and decay amplitudes are as follows:

M_1: u(k_1, -) $\bar{d}$(k_2, +) → W^+(q_1, λ_1) Z(q_2, λ_2),
M_2: W^+(q_1, λ_1) → (p_1, -) $^+{}_1$(p_2, +),
M_3: Z(q_2, λ_2) → $_2$(p_3, h) $^+{}_3$(p_4, -h).

The parameters in parentheses are the particle momentum and helicity, respectively. The phase space of this process has five independent angles defined in the center-of-momentum frame: the production angle (Θ), two polar decay angles ($_1$, $_2$), and two azimuthal decay angles (ϕ_1, ϕ_2), which can be thought of as the rotations of the W^+/Z decay planes (n̂_W, n̂_Z) about the W^+/Z momentum axis, measured relative to the production plane n̂. The three planes are defined as n̂ $k_1 \times q_1$, n̂_W $q_1 \times p_2$, and n̂_Z $p_4 \times q_2$. All these kinematic variables are presented in Fig. 3 [Figure 3: see original paper].

The phase shift from strong dynamics only affects the longitudinal-longitudinal combination of W^+Z modes and shifts the corresponding amplitude by an energy-dependent phase δ. This phase shift enters into the azimuthal angle correlation of the W^+/Z decays as an interference effect between the various polarizations of the vector bosons. To see this, recall that W^+/Z decay produces an azimuthal angular dependence exp(is_zϕ) in the amplitude, where s_z and ϕ are respectively the spin projection and azimuthal angle rotation about the W^+/Z direction. For a given interference term between a general helicity combination (λ_1, λ_2) and the (0,0) combination, we find that the relevant terms in the differential cross-section are of the form:

$$\cos(\lambda_1\phi_1 - \lambda_2\phi_2 + \delta) \supset \sin(\lambda_1\phi_1 - \lambda_2\phi_2)\sin\delta.$$

Thus, the sin($\lambda_1\phi_1$ - $\lambda_2\phi_2$) modes in the azimuthal angle correlation strongly suggest the existence of strong dynamics.

The production amplitudes can be separated by different W^+Z helicity combinations (λ_1, λ_2). Among the nine possible helicity combinations, there are three leading contributions from (-,+), (0,0), and (+,-). The four other significant ones are (-,0) (0,+) and (+,0) (0,-), all of which have a relative suppression m_W/√s compared to the leading modes. The (+,+) and (-,-) combinations are too small to affect the kinematic distributions. We note that: (a) (-,+)/(+,-) dominate as $\Theta \to 0$ or π due to t-channel production; (b) numerically the difference between (-,0) and (0,+) or (+,0) and (0,-) is negligible. We parameterize the production amplitudes as M_1(-,+) = A, M_1(0,0) = Be^{iδ}, M_1(+,-) = C, M_1(-,0) = D, M_1(+,0) = E, M_1(0,+) = F, and M_1(0,-) = G. Here, all amplitudes depend on the center-of-mass energy and on Θ. Similar behavior has been pointed out in $e^+e^- \to W^+W^-$ [**?**].

The full differential cross-section can be obtained from:

$$\frac{d\sigma}{d\cos\Theta d\phi_1 d\phi_2} \propto \left| \sum_{\lambda_1,\lambda_2} M^{(1)}_{\lambda_1\lambda_2}(\Theta) M^{(2)}_{\lambda_1}(\theta_1,\phi_1) M^{(3)}_{\lambda_2}(h,\theta_2,\phi_2) \right|^2,$$

where the W^+, Z decay amplitudes are:

$$M^{(2)}_{\lambda_1}(\theta_1,\phi_1) = g_W|q_1| d_{\lambda_1}(\theta_1) e^{i\lambda_1\phi_1},$$

$$M^{(3)}_{\lambda_2}(\theta_2,\phi_2) = g_Z|q_2| d_{\lambda_2}(\theta_2 + (h-1)\pi/2) e^{-i\lambda_2\phi_2}.$$

Here g_Z^± are the couplings of the Z to different lepton helicities. The polar angle dependent function is d_±( ) = (1±cos)/√2, d_0( ) = sin .

If we integrate over both polar angles $_{1,2}$, there is an approximate cancellation in the sin(ϕ_1+ϕ_2) and sin(ϕ_2) correlations between cos $_2$ > 0 and cos $_2$ < 0 due to g_Z^+ -g_Z^-. Therefore, we only integrate the differential cross-section over either cos $_2$ > 0 or < 0 to obtain:

$$\frac{1}{(2\pi)^2}\frac{d\sigma^{\pm}}{d\cos\Theta d\phi_1 d\phi_2} = \frac{1}{64\pi^2 s}\{H^2 \pm (C^2 - A^2 + G^2 - F^2)$$

$$+B\sin\delta\left[\frac{3\pi}{2}(E-D)\sin\phi_1 + \left(\pm 4(F+G) - \frac{3\pi}{2}(F-G)\right)\sin\phi_2\right.$$

$$+\left(\pm 4(A+C)-\frac{3\pi}{2}(A-C)\right)\sin(\phi_1+\phi_2)\Big]+\cdots\Big\},$$

where = ((g_Z-)2 - (g_Z+)2)/((g_Z-)2 + (g_Z+)2) 0.22 and the ellipsis refers to interference terms with cosϕ-type dependence. The σ^± stands for σ(cos $_2$ 0) and H is the overall background $H^2 = (A^2 + B^2 + C^2 + D^2 + E^2 + F^2 + G^2)$.

In Fig. 4 [Figure 4: see original paper] we plot the relative coefficients of the different sinϕ correlations as a function of cosΘ for s = 1 TeV using the expressions above. Measuring a non-zero coefficient for any of the sinϕ modes is a positive indicator of strong dynamics. Integrating over the entire cosΘ range would lead to cancellations that dilute the significance of the probe. Thus, an optimal strategy is to use a maximum likelihood analysis on the measured dσ^±/d cosΘ dϕ_1 dϕ_2 distribution to look for all non-zero sin($\lambda_1\phi_1$ - $\lambda_2\phi_2$) coefficients.

Procedure

At the LHC, various kinematic ambiguities must be incorporated into the analysis. We choose to study only the sinϕ_1 mode, which yields a much more transparent analysis that is robust to these ambiguities.

We simulate the process pp → W^+Z → $^+$ $^+$ $^-$ with (1 - a^2)F(s) from Eq. (4) multiplying the form factor in the J = 1 channel of the (0,0) helicity amplitude. We take a = 0.8 as a benchmark, assuming such a deviation in Higgs couplings would continue to be allowed with future LHC data. We then scan over different form factor parameters m and Γ, which give rise to different phase shifts. Our event simulation is purely at parton level, which is sufficient for the fully leptonic final state. We use HELAS [?] to calculate the helicity amplitudes for the full process. LHApdf [?] is used to fold in the parton distribution functions for the protons using the PDF set CTEQ6L [?]. An adaptive Monte Carlo package, BASES [?], performs the integration over phase space and studies differential cross-sections. Our simulation approach allows us to insert the form factor specifically into the longitudinal gauge-boson scattering channel.

The cuts used are: (1) ΔR > 0.4 separation between leptons; (2) p_T > 20 GeV and | | < 2.8 cuts on the leptons; (3) Z-reconstruction cut: we require two opposite-sign leptons to reconstruct the Z mass; (4) missing E_T cut > 20 GeV; (5) invariant mass cut of the W/Z system between m ± Γ [?].

When reconstructing events, two misidentification issues lead to a fourfold ambiguity in the kinematics: (a) the u-quark direction is unknown; (b) there is a twofold ambiguity in the neutrino momentum along the beam axis.

First, consider misidentification of the u-quark direction. This leads to Θ → π - Θ, ϕ_1 → π + ϕ_1, ϕ_2 → π + ϕ_2. The azimuthal angle correlation sinϕ_1 is odd under such misidentification. Note that from Fig. 4, the coefficient of sinϕ_1 is also approximately odd under $\Theta \to \pi$ - Θ. Thus, studying the sinϕ_1 mode for either 0.1 < cosΘ < 0.9 or -0.9 < cosΘ < -0.1 makes it robust to misidentifications of the u-quark direction.

The presence of a false solution for the neutrino momentum would distort the azimuthal angle correlations we seek for ϕ_1. However, to study the sinϕ_1 mode, we can simply measure the up-down asymmetry with respect to the $\phi_1 = 0$ (production) plane to find the size of such correlations. We demonstrate that the up-down asymmetry is the same for both true and false solutions.

For a given ϕ_1 azimuthal angle correlation, we have:

$$\left.\frac{d\sigma}{d\phi_1 d\cos\Theta}\right|_{\cos\Theta \gtrless 0} \propto A_0 + A_1 \cos\phi_1 + A_2 \cos 2\phi_1 \pm B_1 \sin\phi_1.$$

We define events going "above" the plane for sinϕ_1 > 0 and "below" the plane for sinϕ_1 < 0. The up-down asymmetry is defined as:

$$AS|_{\cos\Theta \gtrless 0} = \frac{N_+ - N_-}{N_+ + N_-},$$

where N_± are the numbers of up/down events, respectively. The up or down events for ϕ_1 are defined by the sign of the scalar triple product SGN sgn(n̂ · p_2) = sgn((k_1 × q_1) · p_2). For a particular event, if SGN > 0 (< 0) then we increment N_+ (N_-). The normal vector to the production plane n̂ = k_1 × q_1 = k_1 × (p_1 + p_2) is independent of the momentum along the beam direction, and hence SGN is insensitive to the difference between true and false solutions. Additionally, the asymmetry variable has the advantage of being insensitive to several cuts such as rapidity and p_T cuts that would otherwise distort the angular distribution.

Results

If the background fluctuation is Gaussian, the statistical significance of the nonzero asymmetry is given by:

$$\text{Significance} = \frac{|N_+ - N_-|}{\sqrt{N}} = |AS|\sqrt{N},$$

where N = N_+ + N_- is the total number of events. As a rule of thumb, we find that choosing an invariant mass window between m ± Γ optimizes the trade-off between obtaining a large |AS| by being close to the resonance while maintaining a sizeable number of events.

In Table I , we show the cross-section for the process under consideration at the 14 TeV LHC and at a future 50 TeV pp collider for different choices of form factors by varying parameters m and Γ. The form factors we consider lead to typical asymmetries of order 5-10%.

In Table II , we show the significance of the asymmetry measurement at the 14 TeV LHC with 3000 fb^{-1} of data for different form-factor parameters. The results use expected statistics including both W^+Z and W^-Z fully leptonic modes. We find that new wide resonances can be probed at the 3σ level for masses up to more than 900 GeV. This further motivates an extended LHC run should an excess in W±Z be discovered. A future 50 TeV proton-proton collider could probe resonances up to around 1.2 TeV with the same luminosity. An even higher energy at a future collider would make the $\sin(\phi_1+\phi_2)$ interference term dominant and require a different analysis strategy.

Several theoretical and analysis issues could potentially increase these significances in a more sophisticated search. First, in the $\sin\phi_1$ mode search, it is possible to open up the hadronic decay modes of the Z with boosted tagging techniques [?]. Second, using a multivariate analysis and incorporating the $\sin\phi_2$ and $\sin(\phi_1+\phi_2)$ modes could bolster this result. Third, allowing for I = 2 resonances implies that the vector resonance form factor could be scaled by a factor larger than 1 - a^2. Fourth, multiple resonances in a narrow mass window could yield an enhancement in the longitudinal scattering cross-section that would appear as the P(s) factor mentioned earlier.

Conclusions

We have proposed a novel technique to disentangle the dynamics of a strongly coupled EWSB sector by measuring a phase shift in the decay azimuthal angle correlations in massive gauge boson scattering. Our results show that a simple up-down asymmetry in leptons from W decay in pp → W±Z is robust to several event reconstruction ambiguities and serves as a good probe of broad resonances from strong dynamics. This strongly motivates a high-luminosity run of the 14 TeV LHC. A future 50 TeV pp collider could yield conclusive evidence of resonant behavior in the presence of an excess of WZ events at the LHC. Furthermore, we have outlined several analysis strategies and theoretical issues that would significantly increase the reach of searches based on this technique and could lead to a promising signal in the next LHC run.

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[21] See Ref. [5] and references therein.
[22] There could be other non-universal consequences, such as resonances, enhanced di-higgs production rate [6], etc.
[23] See Ref. [7] and references therein.
[24] In practice this requires looking for a broad excess and placing an invariant mass cut on the WZ. We found that our results were not significantly affected by choosing a different invariant mass cut window, indicating that precise optimization of this cut may not be very important.

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Note: Figure translations are in progress. See original paper for figures.