Strongly Interacting W Bosons: Phenomenology of the No-Higgs Case

Wolfgang Kilian*
Institut für Theoretische Teilchenphysik
Universität Karlsruhe, D–76128 Karlsruhe, Germany
(Dated: October 18, 2001)

For a detailed understanding of the physics responsible for electroweak symmetry breaking, one needs access to elastic and inelastic Goldstone scattering amplitudes which are related to the interactions of massive vector bosons. While in a weakly-interacting scenario one expects signatures of Higgs states and their decays in these amplitudes, in the absence of Higgs bosons they become strong in the TeV energy range. A model-independent analysis requires the measurement of those parameters that describe the amplitudes in a low-energy expansion, until new signatures of the underlying physics can be observed directly. This can be done at future hadron and lepton colliders.

I. THE MODEL-INDEPENDENT APPROACH

While the Standard Model with a weakly interacting Higgs sector provides a convenient parameterization for all electroweak precision data, direct searches for Higgs bosons have not been successful so far. This could be caused by the Higgs boson(s) being too heavy to be seen at past and present collider experiments. However, electroweak precision data indicate that, within the context of the Standard Model, the Higgs boson mass is not much larger than 100 GeV [1]. In supersymmetric extensions of the Standard Model in particular, a light Higgs boson is a necessary ingredient. In the LEP 2 runs this mass range has already been explored to some extent, and while there is still a considerable window for a Higgs boson above the current lower mass limit, one should seriously consider the alternative possibility that the light-Higgs scenario is not realized in Nature [2].

In any case, until the degrees of freedom that constitute the Higgs sector have been fully explored and shown to be weakly interacting, the Standard Model with an elementary Higgs boson should be considered as a *model* in the narrow sense. In a purely phenomenological spirit one may step back and use a Higgs-less effective theory as a generic parameterization of the physics of electroweak symmetry breaking [3]. This is known as the *chiral Lagrangian* approach. Only the particles that actually have been observed, vector bosons and fermions, are initially incorporated in the Lagrangian. If there turn out to be additional states, including Higgs or Higgs-like scalars, they may easily be added. Thus, the Standard Model and its extensions are part of the model space covered. In fact, this effective-theory formalism is the most general framework that accounts for spontaneous breaking of the electroweak symmetry.

Formally, the leading terms of the chiral Lagrangian are obtained as the $m_H \to \infty$ limit of the Standard Model. The result is a non-renormalizable quantum field theory. In such a theory, additional interaction terms are needed to regulate higher-order divergences corresponding to unbounded scattering amplitudes. In principle, the number of independent local interactions that have to be introduced would become infinite if one calculated observables to all orders in perturbation theory. In practice, however, the experimental precision is limited, and observables are needed (and calculable) not with infinite accuracy, but typically just to leading (LO) and next-to-leading order (NLO) in the electroweak couplings. Thus, a parameterization in terms of a finite number of parameters is possible and sufficient. The NLO parameters, which cannot be predicted from the low-energy structure of the theory, carry all the information about the Higgs sector that can be obtained without direct observation of new degrees of freedom.

The systematic procedure of starting with the $m_H \to \infty$ limit of the Standard Model, calculating radiative corrections and adding local interactions order by order yields an expansion of scattering amplitudes in terms of

$$\frac{E^2}{\Lambda^2} \quad \text{where } \Lambda = 4\pi v \text{ with } v = (\sqrt{2}G_F)^{-1/2}. \tag{1}$$

Here, E is a linear combination of energies, masses and momenta in a given scattering process. Inserting for E twice the W mass $M_W = gv/2$ one recovers the electroweak loop expansion in terms of $E^2/\Lambda^2 = g^2/16\pi^2$,

*Electronic address: kilian@particle.uni-karlsruhe.de

but in the absence of a light Higgs boson there are other contributions which rise with energy. The scale Λ is associated with the breakdown of perturbation theory. For energies beyond Λ, the low-energy effective theory ceases to be predictive and has to be replaced by another, more fundamental (effective) theory. Consistency in the presence of radiative corrections requires Λ to be of the order $4\pi v \approx 3$ TeV, while if there are new effects below this limit, the scale Λ is lowered accordingly.

The chiral Lagrangian is organized by imposing $SU(2)_L \times U(1)_Y$ invariance on all individual interactions. This symmetry is realized linearly on the fermion fields and on the transversal components of the gauge boson fields. The Higgs sector is generically parameterized by a matrix-valued field Σ with a non-vanishing expectation value v. In a minimal parameterization, it can be written as

$$\Sigma = v \exp \frac{i}{v} w^a \tau^a \tag{2}$$

where w^a, $a = 1, 2, 3$ are three scalar degrees of freedom, the Goldstone bosons. The electroweak symmetry is realized nonlinearly on those fields. In the high-energy limit, their interactions become equivalent with the interactions of longitudinally polarized electroweak gauge bosons [4].

Unless compensated by the presence of Higgs bosons, Goldstone scattering amplitudes rise with energy, and new, independent interaction terms arise at NLO in the perturbative expansion to regulate this effect. They are the only available information about the underlying theory as long as no new degrees of freedom show up in the Higgs sector.

Longitudinally polarized vector bosons couple to other particles proportional to their mass. Thus, the free parameters which depend on the underlying theory of electroweak symmetry breaking enter the amplitudes

- at NLO in the E^2/Λ^2 expansion for $W_L^\pm$ and Z_L interactions;
- at NLO in the loop expansion (energy-independent) for $W_T^\pm$, Z_T and γ interactions;
- at NLO, suppressed by fermion masses, for helicity-changing fermionic interactions;
- at NNLO for helicity-conserving fermionic interactions;
- at NNLO, suppressed by fermion masses, and at NNNLO for pure gluonic interactions.

The leading-order amplitudes do not carry nontrivial information about the Higgs sector. This fact is known as the *low-energy theorem* (LET) [5]. For a meaningful measurement one needs to be sensitive to the NLO, either by projecting onto longitudinally polarized vector bosons at high energies (such that the NLO in E^2/Λ^2 contributes) or via high precision (to be sensitive to loop effects). Unfortunately, the interactions of those particles which are accessible most easily (electrons, light quarks and gluons) are least sensitive, and the effect on transversally polarized vector bosons cannot be enhanced by increasing the collider energy. Therefore, exploring the interactions of Goldstone bosons via the longitudinally polarized states of $W^\pm$ and Z is a key issue for uncovering the mechanism responsible for electroweak symmetry breaking.

II. GOLDSTONE BOSON SCATTERING AMPLITUDES

If the LO part of the chiral Lagrangian is used for computing scattering amplitudes involving Goldstone bosons (resp. longitudinally polarized vector bosons), one obtains expressions that rise with energy and saturate the unitarity bound for partial-wave amplitudes above a certain energy threshold (Fig. 1) [6]. On one hand, this indicates the inconsistency of using only the leading order for extrapolations and the necessity for including higher-order corrections, i.e., the non-renormalizability of the model. On the other hand, one can read off the scale where perturbation theory breaks down, new physics effects come into play, and where one should find hints for the underlying theory of electroweak symmetry breaking. If there is just a Higgs boson in the minimal representation, this scale is replaced by the Higgs mass m_H, and the Higgs resonance restores unitarity to all orders and renders the model weakly interacting up to very high energies. However, the Higgs sector may be more complicated.

As Fig. 1 shows, in quasi-elastic vector boson scattering $WW, ZZ \to WW, ZZ$ the unitarity saturation scale is of the order 1 TeV if higher-order corrections are ignored:

$$\Lambda = \sqrt{8\pi}\, v \approx 1.2 \text{ TeV} \quad (I = J = 0 \text{ partial-wave amplitude}) \tag{3}$$

For fermionic final states, the amplitudes are suppressed by the fermion masses, therefore the corresponding unitarity saturation scales are higher [7]:

$$\Lambda = \frac{8\pi}{\sqrt{3N_c}} \frac{v}{m_f}, \tag{4}$$

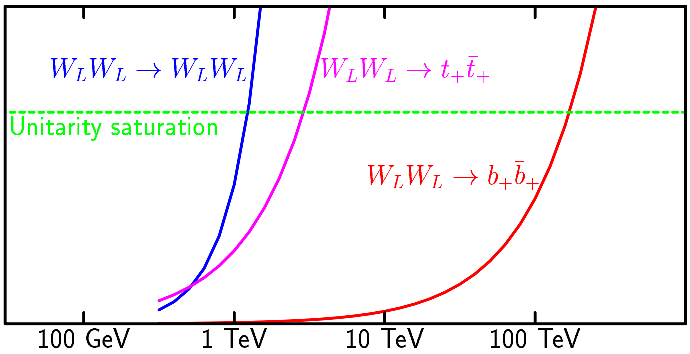


FIG. 1: Energy dependence of various $J = 0$ partial wave amplitudes, calculated to leading order in the chiral Lagrangian approach. The curves correspond to the amplitudes which rise most steeply in the indicated channels, where W_LW_L stands for a linear combination of longitudinally polarized W and Z bosons.

where $N_c = 1$ for leptons and $N_c = 3$ for quarks. (This relation should be taken with a grain of salt since the physics that restores unitarity in quasi-elastic scattering will in general also affect vector boson scattering into fermion pairs.)

Clearly, there must be some new effect in the quasi-elastic vector boson scattering amplitudes that restores unitarity, and if energies in the TeV range are accessible, it will be observable. Depending on the representation of the Higgs sector, a number of scenarios can be imagined:

1. No Higgs: The unitarity bounds are actually saturated, in which case there is a new strong interaction between the Goldstone bosons. This interaction may manifest itself in resonances at and above the TeV scale, and one may think of new constituents which are bound in Goldstone bosons. While this behavior is known from low-energy QCD, quantitatively the situation in the Higgs sector may be much different. For instance, all resonances may be too broad to be resolved, resulting in a strongly interacting continuum of states.

2. A heavy Higgs resonance (m_H = several 100 GeV): This restores unitarity in the first place, but Higgs loop effects make the amplitude rise again, indicating another scale where the model breaks down. Thus, such a Higgs resonance may itself be a bound state in a strongly-interacting theory.

3. A light Higgs particle with Standard Model couplings: Such a particle may in fact be elementary. Quasi-elastic scattering of Goldstone bosons is then expected to be weak at all scales, although there could be surprises. Nevertheless, to get a handle at the Higgs interactions which are associated with electroweak symmetry breaking, one has to look at Goldstone scattering channels again, for instance Higgs pair production $WW/ZZ \to HH$.

4. A light Higgs particle with couplings which significantly deviate from the Standard Model expectation: In that case, unitarity requires additional scalar states (or strong interactions) in the Higgs sector. Depending on the Higgs couplings, the characteristic scale of this new threshold in Goldstone scattering may be significantly higher.

5. If heavy fermions are involved in the physics of electroweak symmetry breaking, the inelastic Goldstone scattering amplitudes $WW/ZZ \to t\bar{t}/b\bar{b}$ should reveal it. If there is a Higgs resonance, these amplitudes will give access to the fermion Yukawa couplings, resolving the issue whether the same Higgs state is associated both with electroweak symmetry breaking and with fermion mass generation.

6. New strong interactions, if they have any resemblance with QCD, may exhibit spontaneous breaking of additional (approximate) chiral symmetries, resulting in new multiplets of pseudo-Goldstone bosons. These will couple to the known Goldstone bosons, and thus be visible in vector boson scattering.

Precision electroweak data are reasonably well described by the Standard Model with a light Higgs boson. While this certainly favorizes weakly interacting scenarios which have a Standard-Model-like *decoupling* limit, the argument is not a strong one. As discussed above, the Higgs sector affects low-energy observables only indirectly. At present, the only meaningful information about the Higgs sector is contained in two parameters S and T which, in the Standard Model context, depend logarithmically on the Higgs mass [8]. The success of the Standard Model prediction for these parameters is remarkable, but it is still quite possible that this is a coincidence [9]. The information contained in Goldstone scattering amplitudes (in particular, the presence or absence of Higgs resonances, of course) will resolve this issue.

Clearly, Goldstone (resp. longitudinal vector boson) scattering amplitudes are an important source of information about the Higgs sector in any scenario, weakly interacting or strongly interacting. Since the former is covered in detail in other contributions to this Snowmass workshop, in the remainder of this report we will concentrate on the latter case.

There are three independent amplitudes involved in quasi-elastic vector boson scattering: $WWWW$, $WWZZ$ and $ZZZZ$. At NLO in the chiral expansion, the amplitudes depend on five new parameters. If one assumes that a custodial $SU(2)_C$ symmetry [10] is realized in the Higgs sector, only two of them are independent which may be taken as α_4 and α_5 [3]:

$$\mathcal{L}_4 = \alpha_4 (\mathrm{tr}\,[V_\mu V_\nu])^2 \tag{5}$$

$$\mathcal{L}_5 = \alpha_5 (\mathrm{tr}\,[V_\mu V^\mu])^2 \tag{6}$$

$$\tag{7}$$

where, in unitary gauge,

$$V_\mu = i\frac{g}{\sqrt{2}}(W_\mu^+ \tau^+ + W_\mu^- \tau^-) + i\frac{g}{2c_w} Z_\mu \tau^3. \tag{8}$$

The quasi-elastic vector boson scattering amplitudes can be expressed in terms of a single function $A(s,t,u)$,

$$A(W_L^- W_L^- \to W_L^- W_L^-) = A(t,s,u) + A(u,t,s) \tag{9}$$

$$A(W_L^+ W_L^- \to W_L^+ W_L^-) = A(s,t,u) + A(t,s,u) \tag{10}$$

$$A(W_L^+ W_L^- \to Z_L Z_L) = A(s,t,u) \tag{11}$$

$$A(W_L^- Z_L \to W_L^- Z_L) = A(t,s,u) \tag{12}$$

$$A(Z_L Z_L \to Z_L Z_L) = A(s,t,u) + A(t,s,u) + A(u,t,s) \tag{13}$$

where, to NLO in the low-energy expansion and including leading radiative corrections,

$$\begin{aligned}\mathrm{Re}\,A(s,t,u) = \frac{s}{v^2} + \frac{1}{16\pi^2 v^4}\left\{ -\frac{(t-u)}{6}\left[t \ln\frac{-t}{\mu^2} - u \ln\frac{-u}{\mu^2}\right] - \frac{s^2}{2}\ln\frac{s}{\mu^2}\right\} \\ + \alpha_4 \frac{4(t^2+u^2)}{v^4} + \alpha_5 \frac{8s^2}{v^4}.\end{aligned} \tag{14}$$

Without $SU(2)_C$ symmetry, the amplitudes involve three more independent interaction terms,

$$\mathcal{L}_6 = \alpha_6\, \mathrm{tr}\,[V_\mu V_\nu]\,\mathrm{tr}\,[TV^\mu]\,\mathrm{tr}\,[TV^\nu] \tag{15}$$

$$\mathcal{L}_7 = \alpha_7\, \mathrm{tr}\,[V_\mu V^\mu]\,\mathrm{tr}\,[TV_\nu]\,\mathrm{tr}\,[TV^\nu] \tag{16}$$

$$\mathcal{L}_{10} = \tfrac{1}{2}\alpha_{10} (\mathrm{tr}\,[TV_\mu]\,\mathrm{tr}\,[TV_\nu])^2 \tag{17}$$

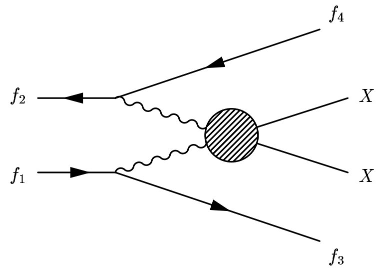
and the scattering amplitudes are modified accordingly.

III. PHENOMENOLOGY

Scattering amplitudes of longitudinally polarized vector bosons can be accessed in collider experiments in various ways:

1. At high energies, electroweak gauge bosons can be emitted from fermions, thus providing direct access to the amplitudes in question: The processes

$$f_1 f_2 \to f_3 f_4 X X \tag{18}$$



contain $VV \to XX$ as a subprocess, where $V = W, Z$, and XX is any of the final states discussed above: vector bosons, Higgs bosons, heavy fermions, pseudo-Goldstone bosons, etc.

In the energy range where the mass of the intermediate vector bosons can be neglected, their virtuality is typically negligible as well, such that they can be treated as partons inside the incoming fermions [11]. At large transverse momenta, the emission of longitudinally polarized vector bosons is suppressed by the incoming fermion mass, so in the interesting range where the processes are sensitive to details of the Higgs sector the transverse momentum of the outgoing fermions is small, of order $M_W/2$. On the other hand, the final state particles X will in general be observed at large scattering angle.

Approximately, the total cross section of a process of the type (18) is given by

$$\sigma(s) = \sum_{\lambda_1\lambda_2} \int dx_1\, dx_2\, F_{\lambda_1}(x_1)\, F_{\lambda_2}(x_2)\, \hat{\sigma}_{\lambda_1\lambda_2}(x_1 x_2 s) \tag{19}$$

where $\hat{\sigma}$ is the on-shell cross section of the subprocess $VV \to XX$, and one has to sum over the vector boson polarization states $\lambda = \pm, 0$. The vector boson structure functions read

$$F_{W,\pm}(x) = \frac{g^2}{8\pi}\,\frac{1+\bar{x}^2}{x}\,\ln\frac{p_\perp^2}{M_W^2} \tag{20}$$

$$F_{Z,\pm}(x) = \frac{g^2}{8\pi c_w^2}\left[(t_3^f - 2q_f s_w^2)^2 + (t_3^f)^2\right]\frac{1+\bar{x}^2}{x}\,\ln\frac{p_\perp^2}{M_Z^2} \tag{21}$$

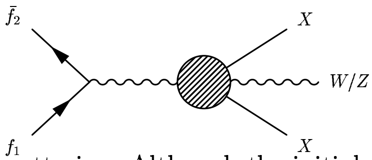
$$F_{W,0}(x) = \frac{g^2}{4\pi}\,\frac{\bar{x}}{x} \tag{22}$$

$$F_{Z,0}(x) = \frac{g^2}{4\pi c_w^2}\left[(t_3^f - 2q_f s_w^2)^2 + (t_3^f)^2\right]\frac{\bar{x}}{x} \tag{23}$$

where $\bar{x} \equiv 1 - x$, $p_\perp$ denotes a cutoff on the transverse momentum, g is the $SU(2)_L$ gauge coupling, c_w and s_w are the cosine and sine of the weak mixing angle, respectively, t_3^f is the $SU(2)_L$ quantum number, and q_f is the normalized electric charge of the incoming fermion. Depending on the detailed energy dependence of the subprocess cross section $\hat{\sigma}$, the total cross section (19) rises with energy at least logarithmically. Unfortunately, part of the total available energy goes into the final-state fermions, and the effective scale where the subprocess amplitude is probed is just a fraction of the collider c.m. energy.

2. Associated production

$$f_1 \bar{f}_2 \to VXX \tag{24}$$

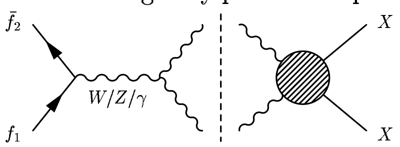


is sensitive to the same amplitudes as vector boson scattering. Although the initial-state fermions do not couple to longitudinally polarized vector bosons directly, the large virtuality of the s-channel Z or W boson induces a mixing of transversal and longitudinal states. As far as the Higgs sector is involved in this scattering, the final-state vector boson V is also predominantly longitudinally polarized. XX may denote a vector boson pair (triple vector boson production), the decay products of a Higgs resonance (Higgs-strahlung), a Higgs pair, or any other state which has a significant coupling to the symmetry-breaking sector.

The cross section for this class of processes asymptotically falls off like $1/s$. This certainly holds for weakly interacting particles, but even in the strongly-interacting scenario where Goldstone scattering amplitudes initially rise proportional to s, above the unitarity saturation threshold the amplitude is bounded and a $1/s$ falloff sets in. Thus, this type of processes is useful at low collider energies, not too much above the production threshold of the VXX state. At higher energies, the cross section for vector boson fusion is more important. However, depending on the process under consideration this asymptotic region may be beyond the reach of collider experiments, such that associated production provides a better access to the interesting observables.

3. Rescattering corrections, e.g. in vector boson pair production, are also sensitive to vector boson scattering. A vector boson $2 \to 2$ scattering amplitude provides the imaginary part of the pair production amplitude:

$$f_1 \bar{f}_2 \to XX \tag{25}$$



This is a NLO effect, and if one is looking for the NLO parameters that carry nontrivial information about the Higgs sector, they only make up a NNLO contribution to the pair production process which is difficult to extract without ambiguities. However, the situation is more favorable if there is a strong resonance in the Goldstone boson scattering amplitude. The rescattering effect takes place at the full collider energy. The fact that the $J = 1$ part of the scattering amplitude is projected onto in this process makes it uniquely sensitive to vector states, and with sufficient precision their presence can be detected significantly beyond the kinematical reach of the collider [12].

IV. EXPERIMENTS AT HADRON COLLIDERS

With sufficient c.m. energy, hadron colliders are well suited for the measurement of Goldstone scattering amplitudes in the vector boson fusion channels (18). In addition, associated production channels (24) can be exploited, where the recoiling vector boson serve as a tag for identifying the signal above the background. Unfortunately, the large hadronic background rate makes the identification of light Higgs bosons, which predominantly decay hadronically, very difficult. A strongly interacting Higgs sector (with or without a heavy Higgs boson) is somewhat easier to study since there are vector boson final states which have a significant leptonic decay branching fraction.

Here, our main focus is on the no-Higgs case. Quasi-elastic scattering of vector bosons (18) is observable up to (roughly) two thirds of the partonic c.m. energy, which is typically less than one tenth of the available collider energy. For higher energies, the rate quickly drops, so at LHC one expects no significant rate of vector boson scattering processes beyond about 1 TeV for 14 TeV total collider energy. Even in strongly interacting scenarios it is not guaranteed that resonances fall into this energy range (if there are any), so it is conceivable that experiments will have to concentrate on precision measurements of the NLO parameters which describe the behavior of quasi-elastic vector boson scattering. While at the Tevatron the collider energy is too low for a meaningful measurement of vector boson scattering, a VLHC with 40 TeV c.m. energy can essentially cover the interesting energy range up to $4\pi v \approx 3$ TeV [13].

The strategy for isolating the quasi-elastic scattering signal in hadronic collisions is similar to the identification of a heavy Higgs boson in its vector boson decay modes. One should tag the forward-scattered quarks [cf. (18)] as energetic jets with low transverse momentum. For typical values $p_T = M_W/2$ and $E \approx 1$ TeV, they appear at $\eta \approx \pm 4$. Concentrating on leptonic vector boson decays, one will also apply a jet veto in the central region which significantly reduces the hadronic background. The scattering processes that can be observed are

$$q_1 q_2 \to q_3 q_4 W^+ W^- \tag{26}$$

$$q_1 q_2 \to q_3 q_4 ZZ \tag{27}$$

$$q_1 q_2 \to q_3 q_4 W^\pm Z \tag{28}$$

$$q_1 q_2 \to q_3 q_4 W^- W^-, \;\; q_3 q_4 W^+ W^+ \tag{29}$$

While the first two channels are identical to the ones one is looking for in heavy Higgs production, enhancements in the other channels point into a different direction. The WZ final state, for instance, is typical for the decay of a vector resonance. In the non-resonant case, one will combine all channels for a measurement of the NLO chiral parameters.

In any case, after accounting for reducible backgrounds (e.g., WW pair production, top quark production, QCD processes), in isolating the strong scattering signal one has to cope with the large production rate of transversally polarized vector bosons. To project onto longitudinal polarization states, one would like to reduce the contribution of low vector boson pair invariant mass in the event sample and to exploit angular distributions in the vector boson decays. Due to the limited capabilities of Monte Carlo generation packages this has only partly been possible in the published studies. However, the results show that at the LHC there is significant sensitivity to the parameters of interest (Fig. 2).

V. EXPERIMENTS AT LEPTON COLLIDERS

The next generation of e^+e^- colliders that is currently proposed opens an opportunity for precision measurements that directly address the physics of electroweak symmetry breaking. The low rate of hadronic background in lepton machines allows to make use of all important decay channels of particles associated with the Higgs sector. Furthermore, the accurate knowledge of the c.m. energy in lepton collider events allows for deriving the missing invariant mass, which is crucial for the analysis of final states containing neutrinos. Nevertheless, the low signal rates of many of the processes of interest require a high luminosity.

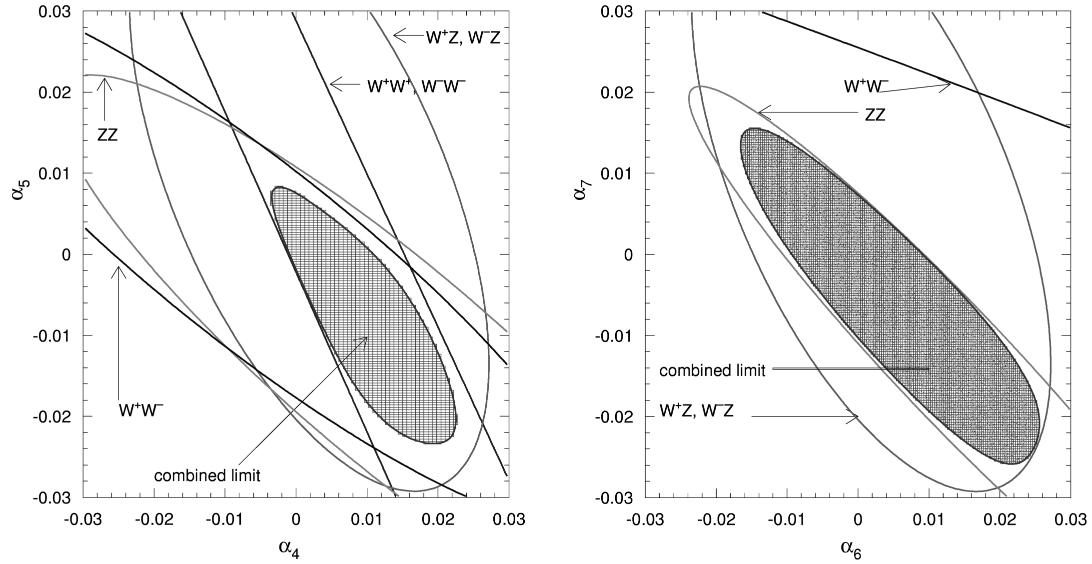


FIG. 2: 1-σ exclusion regions in the $\alpha_{4,5}$ (left) and $\alpha_{6,7}$ (right) planes which results from a cut-based analysis of quasi-elastic vector boson scattering processes. The contours are based on the hypothesis that the actual value of all parameters is zero, assuming an integrated luminosity of 100 fb^{-1} at LHC [14].

The first stage of linear colliders is foreseen to reach a c.m. energy of 500 GeV. At this energy, in the Higgs-less case there is only a marginal possibility to measure vector boson scattering amplitudes in the fusion channels. However, associated production of W and Z pairs, i.e., triple vector boson production

$$e^+e^- \to W^+W^-Z,\ ZZZ \tag{30}$$

is accessible, and there is a certain sensitivity on the NLO parameters $\alpha_{4,5,6,7,10}$ in these processes. Studies show that indeed meaningful constraints on these parameters can be set, adding complementary information to the results available from LHC analyses [15]. Of course, in the light-Higgs case the situation is much more favorable at an e^+e^- collider: Analyzing the associated-production channels of single and double Higgs-strahlung in all decay modes, one may be able to draw an accurate picture of the Higgs sector even without the ability to access higher energies.

When, in a second stage, the c.m. energy of e^+e^- collisions will be increased up to $800\ldots1000$ GeV, the achievable sensitivity in the measurement of quasi-elastic vector boson scattering becomes more than comparable to the LHC reach, depending on the luminosity that can be accumulated. Polarizing the electron and positron beams is equivalent to an additional increase of up to a factor of 4 in effective signal luminosity accompanied by only a minor enhancement of the important backgrounds.

While the topology of the signal Feynman diagrams is equivalent to those relevant for hadron collider processes (18), the experimental signature looks quite different. Instead of forward jets, in e^+e^- collisions one has to deal with forward neutrinos, i.e., a large missing invariant mass in the signal events. The missing transverse momentum is limited, of the order $M_W/2$. One has to remove events initiated by photons radiated from the incoming electron or positron, so to veto forward electrons or positrons a good coverage of the forward detector region is essential. In the central detector region one has to look for the decay products of the final-state W and Z bosons which may be either leptons or jets. The latter decays are favored because additional neutrinos in leptonic W and invisible Z decays contribute to the overall missing energy-momentum.

The channels that are easily accessible are [16]

$$e^+e^- \to \bar\nu_e\nu_e W^+W^- \tag{31}$$
$$e^+e^- \to \bar\nu_e\nu_e ZZ \tag{32}$$
$$e^-e^- \to \nu_e\nu_e W^-W^- \tag{33}$$

where the last process requires running of the collider in the e^-e^- mode [17]. Furthermore, one is interested in

$$e^+e^- \to e^\pm\nu_e W^\mp Z \tag{34}$$
$$e^+e^- \to e^+e^- W^+W^- \tag{35}$$
$$e^+e^- \to e^+e^- ZZ \tag{36}$$

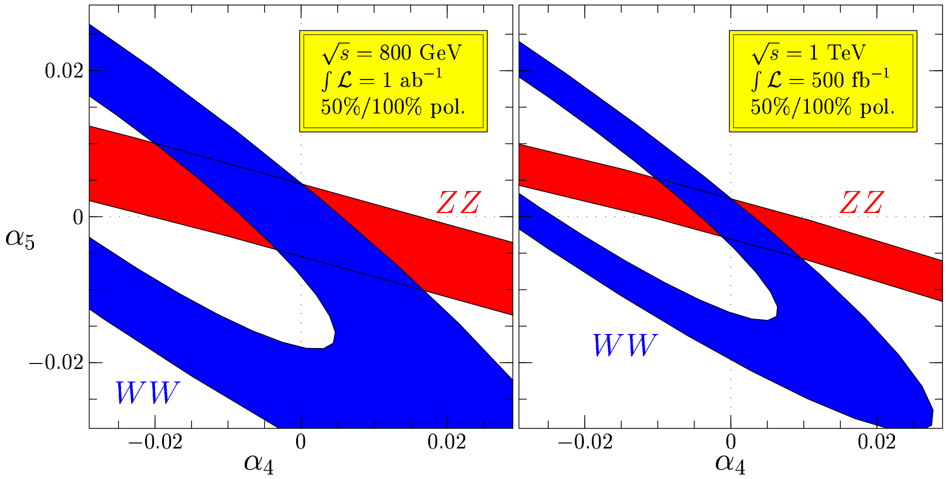


FIG. 3: Allowed regions in the α_4-α_5 plane resulting from a cut-based analysis of quasi-elastic vector boson scattering in polarized e^+e^- collisions for two different sets of collider parameters. The exclusion contours are based on the hypothesis that the actual values are $\alpha_4 = \alpha_5 = 0$ [18].

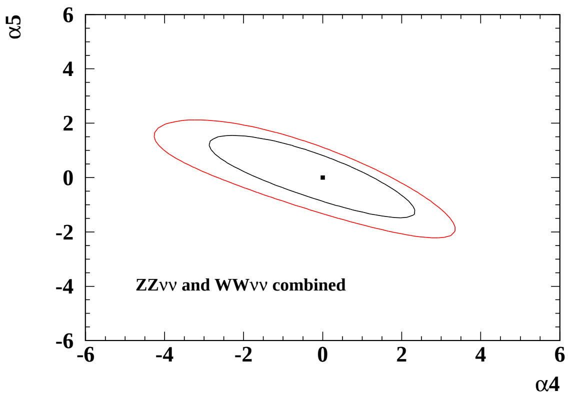


FIG. 4: Allowed regions in the α_4-α_5 plane resulting from a likelihood analysis of simulated six-fermion event samples corresponding to quasi-elastic vector boson scattering in polarized e^+e^- collisions [19]. The exclusion contour is based on the hypothesis that the actual values are $\alpha_4 = \alpha_5 = 0$. The assumed integrated luminosity is $\int \mathcal{L} = 1$ ab^{-1} at an energy of $\sqrt{s} = 800$ GeV with 80 % (40 %) electron (positron) polarization. The inner and outer contours represent 68 % C.L. and 90 % C.L. limits, respectively. Note that a factor of $16\pi^2$ is absorbed in the definition of α_4 and α_5; to compare the results with Fig. 3, for instance, the numbers on the axes have to be multiplied by $1/16\pi^2 = 0.00633$.

but these processes have a lower signal rate and are plagued by larger background, so they are of lesser importance.

Quantitatively, a cut-based analysis of the total cross sections in the W^+W^- and ZZ channels yields a sensitivity to the NLO chiral parameters $\alpha_{4,5}$ which is in the range $1/16\pi^2 \sim 0.01$ one is interested in (Fig. 3). High luminosity is needed, and a good separation of the dijet invariant mass in hadronic W and Z decays must be achieved to arrive at these results. The shown two-parameter plots implicitly assume $SU(2)_C$ symmetry. To get information on $SU(2)_C$-violating couplings, one should analyze the angular distributions in detail, add more channels [the e^-e^- mode (33) and the difficult channels (34–36)], and combine the results with the LHC data.

A better sensitivity estimate can be obtained from a complete simulation of the six-fermion final state in the processes (31, 32). The angular distributions of the final-state fermions can be used for projecting onto longitudinally polarized W and Z bosons, and the sample can be studied event by event in a likelihood analysis, taking into account the expected detector effects. The results show that it is actually possible to further improve the constraints obtained from a phenomenological cut-based analysis, resulting in the exclusion contour shown in Fig. 4.

The possibility that the heavy top quark plays an important role in the physics of electroweak symmetry breaking [20] opens interesting prospects for inelastic Goldstone scattering amplitudes with heavy-quark final states. Unfortunately, they are even more difficult to access than quasi-elastic scattering amplitudes. Preliminary studies show that while at an effective energy below 1 TeV the amplitudes can probably be measured only if they exhibit resonant behavior, with more c.m. energy available there is a sensitivity on the contact interactions of Goldstone bosons and heavy quarks, an important new source of information [21].

VI. CONCLUSIONS

To study the physics of electroweak symmetry breaking in detail, one has to precisely measure the scattering amplitudes which involve Goldstone bosons, equivalent to the scattering amplitudes of longitudinally polarized vector bosons. This is true in any scenario of electroweak symmetry breaking. In particular, if there is no Higgs boson (or if the Higgs resonance is heavy), these scattering amplitudes become strong in the TeV range. Detailed information about their structure will then be an essential ingredient in our knowledge about the physics of mass generation.

The next generation of colliders, LHC and e^+e^- linear colliders together, will be needed to achieve this goal. Coverage of the energy range up to about 1 TeV and the ability to precisely measure cross sections and distributions in all channels is essential here. Nevertheless, while the physics of electroweak symmetry breaking and mass generation will have become much clearer after these measurements have been performed, it is conceivable that the origin of these effects cannot be fully uncovered without access to the multi-TeV range. The VLHC, CLIC and muon collider projects address this issue, and the ability of directly producing new resonances and states significantly above 1 TeV may become important for a deeper understanding of the mechanism responsible for electroweak symmetry breaking.

Acknowledgments

This work was supported by the DFG Graduiertenkolleg “Elementarteilchenphysik an Beschleunigern” and by the DFG Forschergruppe “Quantenfeldtheorie, Computeralgebra und Monte-Carlo-Simulation”.

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[1] M. Swartz, these proceedings; J. Bagger, these proceedings.

[2] R.S. Chivukula, these proceedings; E.H. Simmons, these proceedings.

[3] T. Appelquist and C. Bernard, Phys. Rev. **D22**, 200 (1980); A. Longhitano, Phys. Rev. **D22**, 1166 (1980); Nucl. Phys. **B188**, 118 (1981).

[4] C.E. Vayonakis, Lett. Nuovo Cim. **17**, 383 (1976); M.S. Chanowitz and M.K. Gaillard, Nucl. Phys. **B261**, 379 (1985); G.J. Gounaris, R. Kögerler, and H. Neufeld, Phys. Rev. **D34**, 3257 (1986); Y.-P. Yao and C.-P. Yuan, Phys. Rev. **D38,** 2237 (1988); J. Bagger and C. Schmidt, Phys. Rev. **D34**, 264 (1990).

[5] S. Weinberg, Phys. Rev. Lett. **17**, 616 (1966); M.S. Chanowitz, M. Golden, and H. Georgi, Phys. Rev. **D36**, 1490 (1987).

[6] B. Lee, C. Quigg, and H. Thacker, Phys. Rev. Lett. **38**, 883 (1977); Phys. Rev. **D16**, 1519 (1977); D. Dicus and V. Mathur, Phys. Rev. **D7**, 3111 (1973).

[7] T. Appelquist and M.S. Chanowitz, Phys. Rev. Lett. **59**, 2405 (1987); (E) **60**, 1589 (1988); F. Maltoni, J.M. Niczyporuk, and S. Willenbrock, hep-ph/0106281; S. Willenbrock, these proceedings.

[8] M. Peskin and T. Takeuchi, Phys. Rev. Lett. **65**, 964 (1990); Phys. Rev. **D46**, 381 (1992).

[9] M.E. Peskin and J.D. Wells, Phys. Rev. **D64**, 093003 (2001); M.E. Peskin, these proceedings.

[10] M. Veltman, Act. Phys. Pol. **B8**, 475 (1977); Nucl. Phys. **B123**, 89 (1977); P. Sikivie, L. Susskind, M. Voloshin, and V. Zakharov, Nucl. Phys. **B173**, 189 (1980).

[11] M.S. Chanowitz and M.K. Gaillard, Phys. Lett. **B142**, 85 (1984); G.L. Kane, W.W. Repko, and W.R. Rolnick, Phys. Lett. **B148**, 367 (1984); S. Dawson, Nucl. Phys. **B249**, 42 (1985); J. Lindfors, Z. Phys. **C28**, 427 (1985); J.F. Gunion, J. Kalinowski, and A. Tofighi-Niaki, Phys. Rev. Lett. **57**, 2351 (1986).

[12] T. Barklow, in: A. Para and H.E. Fisk (eds.), *Proceedings of the 5th Int. Linear Collider Workshop (LCWS 2000)*, Am. Inst. of Physics, 2001.

[13] J. Bagger, V. Barger, K. Cheung, J. Gunion, T. Han, G.A. Ladinsky, R. Rosenfeld, and C.-P. Yuan, Phys. Rev. **D49**, 1246 (1994); Phys. Rev. **D52**, 3878 (1995).

[14] A.S. Belyaev, O.J.P. Eboli, M.C. Gonzalez-Garcia, J.K. Mizukoshi, S.F. Novaes, and I. Zacharov, Phys. Rev. **D59**, 015022 (1999).

[15] V. Barger, T. Han, and R.J.N. Phillips, Phys. Rev. **D39**, 146 (1989); S. Dawson, A. Likhoded, G. Valencia, and O. Yushchenko, in: *Snowmass 1996, New directions for high-energy physics*, SLAC 1997; O. Eboli, M. Gonzalez-Garcia, and J. Mizukoshi Phys. Rev. **D58**, 034008 (1998).

[16] V. Barger, K. Cheung, T. Han, and R.J.N. Phillips, Phys. Rev. **D52**, 3815 (1995).
[17] T. Han, Int. J. Mod. Phys. **A11**, 1541 (1996); W. Kilian, Int. J. Mod. Phys. **A15**, 2387 (2000); W. Kilian, these proceedings.
[18] E. Boos, H.-J. He, W. Kilian, A. Pukhov, C.-P. Yuan, and P.M. Zerwas, Phys. Rev. **D57**, 1553 (1998); Phys. Rev. **D61**, 077901 (2000).
[19] R. Chierici, S. Rosati, and M. Kobel, LC-PHSM-2001-038.
[20] B. Dobrescu, these proceedings.
[21] F. Larios, T. Tait, and C.-P. Yuan, Phys. Rev. **D57**, 3106 (1998); T. Han, Y.J. Kim, A. Likhoded, and G. Valencia, Nucl. Phys. **B593**, 415 (2001); J. Alcaraz and E. Ruiz Morales, Phys. Rev. Lett. **86**, 3726 (2001).