PHYSICS OF THE W BOSON AT FUTURE LINEAR COLLIDERS* **

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A survey is given of the various aspects of W-boson physics at the next generation of linear colliders. In particular, it is indicated how the W boson can help us improve our understanding of the mechanism of mass generation and the structure of the non-abelian gauge-boson interactions. Also the topics of radiative corrections and gauge invariance are briefly addressed.

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1. Introduction

During the last five years the high-energy physics community has investigated the viability of a new high-energy linear e^+e^- collider [1]. According to the various designs for this so-called Next Linear Collider (NLC), an energy in the range between 500 and 1500 GeV seems feasible. By using the back-scattered laser-beam technique [2], it is possible to convert the $e^\pm$ beams into photon beams with comparable energy and luminosity. In this way the NLC could be operated as a e^+e^-, e^-e^-, $e\gamma$, or $\gamma\gamma$ collider. Such a versatile machine should prove an excellent tool in our quest to understand nature. It will provide stringent tests of the Standard Model (SM) of electroweak interactions [3], making it highly sensitive to signals of new physics. If any physics beyond the SM exists, it will reveal itself in the production of new particles (direct signals) or in deviations in the interactions between the SM particles (indirect signals).

This survey is dedicated to the role played by the W boson. In this context the most important issues to be addressed at the NLC are the investigation of the triple and quartic gauge-boson couplings, and a detailed

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study of the symmetry-breaking sector. This should shed some light on two important outstanding problems in present-day high-energy physics.

The first question concerns the nature of the non-abelian interactions between the electroweak gauge bosons. Experiments at the Tevatron and LEP1 have provided us with the first direct [4] and indirect [5] evidence for the existence of such interactions. The results are, however, far from conclusive, since $\mathcal{O}(1)$ deviations from the SM couplings are still allowed. At future high-energy collider experiments the sensitivity to these non-abelian gauge-boson couplings will be increased significantly. This will either allow a verification of the SM couplings at the permille level or open a window to physics beyond the SM.

The second question that should be addressed at the NLC concerns the mechanism of electroweak symmetry breaking. Are the longitudinal weak gauge-boson modes indeed generated by means of the SM Higgs mechanism, or has nature chosen another option? In this context there are two scenarios for the electroweak symmetry-breaking sector. In one scenario the theory of the fundamental interactions is assumed to be applicable to very high energies, *i.e.* up to the grand-unification scale or Planck mass (10^{16}–10^{19} GeV). In such a scenario (*e.g.* supersymmetry) the symmetry-breaking sector consists of one or more elementary (pseudo-)scalar Higgs fields, which are weakly coupled at low energies. The mass of the lowest-lying Higgs state is predicted to be relatively small, *i.e.* below 200 GeV [6]. The second scenario involves the possibility of having new strong interactions at TeV energies, related to the mechanism of mass generation. Such a scenario in general excludes the presence of a low-lying Higgs state; in fact, Higgs-like states might be absent altogether. Whatever the underlying theory of the strong interactions may be, these strong interactions will manifest themselves in the form of strongly-interacting longitudinal gauge bosons. After all, these longitudinal gauge-boson modes are a direct consequence of the mechanism of mass generation. This is reminiscent of the role played by the pions in hadron dynamics. In this survey the emphasis will be on the strongly-interacting scenario, since it entails the absence of direct signals at low energies. From the viewpoint of W-boson physics this is the most interesting situation.

In order to successfully achieve the physics goals at the NLC, a very accurate knowledge of the SM predictions for the various observables is mandatory. It has no use trying to perform high-precision tests of non-abelian gauge-boson couplings and strongly-interacting longitudinal gauge-boson interactions when the SM predictions do not have a matching precision. To this end a critical assessment is given as to what SM ingredients are required in this respect. This involves a proper understanding of radiative corrections as well as a proper treatment of finite-width effects. The weak gauge bosons are unstable particles and experience has learned us that gauge invariance is in jeopardy when it comes to including the finite widths of these particles. Needless to say that this can have large repercussions on the reliability of the SM predictions.

The outline of this survey is as follows. In Section 2 a discussion is given of the process of longitudinal gauge-boson scattering and its intimate relation to the symmetry-breaking sector. In Section 3 the topic of physics beyond the SM is addressed in a more systematic way, using the concept of effective Lagrangians and anomalous couplings. In Section 4, finally, the SM predictions are considered, with emphasis on the issue of gauge invariance.

2. Longitudinal gauge-boson scattering

As mentioned before we will assume the absence of a low-lying Higgs state or any alternative thereof, excluding the presence of direct signals of these states at sub-TeV energies. The most sensitive probe of the symmetry-breaking sector will in that case be longitudinal gauge-boson scattering, since the longitudinal gauge-boson modes are a direct consequence of the mechanism of mass generation.

2.1. Strong interactions between longitudinal gauge bosons in the SM

In the SM the longitudinal gauge-boson modes are supplied by the would-be Nambu–Goldstone bosons, leaving behind just one elementary scalar Higgs field (H). This is achieved by the spontaneous breakdown of the $SU(2)_L \times U(1)_Y$ symmetry through the Higgs mechanism. In order to reveal the distinctive nature of the longitudinal gauge-boson modes, high energies are required. At rest longitudinal (L) and transverse (T) modes are related by mere rotations, but at high energies they are quite different, since only the longitudinal modes are affected by boosts in the direction of flight. A simple investigation of the polarization vector $\varepsilon^\mu(k)$ of a massive gauge boson with momentum $k^\mu = (E, \vec{k})$, mass M, and velocity $\beta = \sqrt{1 - M^2/E^2}$ reveals:

$$\begin{aligned} \varepsilon_T^\mu(k) &= (0, \vec{\epsilon}) \quad \text{with} \quad \vec{\epsilon}\cdot\vec{k} = 0 \;\; \text{and} \;\; \vec{\epsilon}^{\,2} = 1 \, , \\ \varepsilon_L^\mu(k) &= \frac{k^\mu}{\beta M} - \frac{M}{\beta E}\,(1, \vec{0}\,) \equiv \frac{k^\mu}{\beta M} + V^\mu \, . \end{aligned} \tag{1}$$

From this it should be clear that any amplitude involving longitudinal gauge bosons has the tendency to diverge at high energies as a result of factors $k/M \sim E/M$. In gauge theories, however, gauge cancellations take place, resulting in properly behaved cross-sections for all energies, *i.e.* cross-sections that do not grow as a positive power of E. This is reflected by the fact that the leading-energy term k^μ/M in (1) can be related by means of Ward identities to the corresponding would-be Goldstone mode, which is not subject to gauge cancellations.

Consider now the process $W_L^+W_L^- \to Z_L Z_L$ in the limit $M^2_{W,Z} \ll (E^2, M_H^2)$, where E stands for the energy of the particles in the centre-of-mass system. At this point we can make use of the equivalence theorem [7], which states that for $E^2 \gg M^2_{W,Z}$ the amplitudes for longitudinal gauge-boson scattering are in leading-energy approximation equivalent to the amplitudes for the corresponding would-be Goldstone bosons. The terms involving the remnant V^μ occurring in (1) are suppressed by powers of $M_{W,Z}/E$[1]. Hence, it suffices to study the process $\phi^+\phi^- \to \chi\chi$ (see Fig. 1), with $\phi^\pm$ and χ the would-be Goldstone bosons responsible for generating the masses of the $W^\pm$ and Z gauge bosons, respectively. This process is not subject to gauge cancellations and is therefore easier to handle. The lowest-order matrix element is given by

$$\mathcal{M}(W_L^+W_L^- \to Z_L Z_L) \approx \mathcal{M}(\phi^+\phi^- \to \chi\chi) \approx \frac{-sM_H^2}{v^2(s-M_H^2)}\,, \tag{2}$$

where $\sqrt{s} = 2E$ stands for the total centre-of-mass energy and $v = 246$ GeV for the electroweak symmetry-breaking scale.

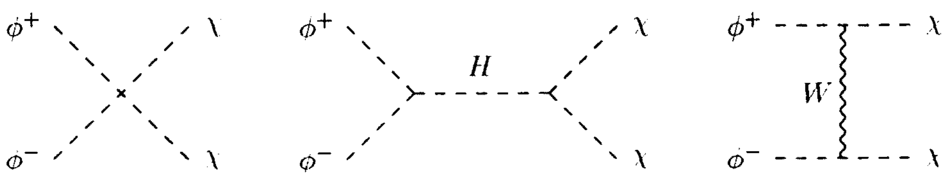

Fig. 1. The Feynman diagrams for the lowest-order process $\phi^+\phi^- \to \chi\chi$. In the limit $M^2_{W,Z} \ll (E^2, M_H^2)$ the third diagram is suppressed with respect to the other two.

If the Higgs boson is very heavy ($E^2 \ll M_H^2$) an interesting phenomenon occurs. Naively the amplitude for the process $W_L^+W_L^- \to Z_L Z_L$ is expected to diverge like E^4, but after the gauge cancellations have taken place an E^2 behaviour remains. To be more precise: $\mathcal{M} \to s/v^2$. This behaviour is purely a consequence of the fact that the symmetry is broken and does

[1] Beyond lowest-order level proper care has to be taken with external self-energies and renormalization factors.

not contain any information on the dynamics responsible for the symmetry breaking, *i.e.* the Higgs. It is prescribed by the Low-Energy Theorem (LET) [8] corresponding to the broken symmetry, which states that the would-be Goldstone bosons decouple at low energies (up to gauge and Yukawa couplings). We will come back to this point later on. Obviously something must happen in the TeV range in order to salvage unitarity, after all the SM is unitary by construction; in other words, strong-interaction effects are expected in that regime. Indeed, in the heavy-Higgs approximation the quartic Higgs self couplings $\lambda \propto M_H^2/v^2$ are large and higher-order corrections $\propto M_H^2/(4\pi v)^2$ can not be discarded. Two natural scales govern the dynamics of these strong-interaction effects: M_H (resonance) and $4\pi v \sim 3$ TeV (corrections). The Goldstone bosons, which decouple at low energies, will start to interact strongly if $E = \mathcal{O}(M_H, 4\pi v)$. As a result of the equivalence theorem, the same holds for the longitudinal gauge bosons.

Recapitulating: the LET behaviour of longitudinal gauge-boson scattering at intermediate energies ($M_{W,Z}^2 \ll E^2 \ll M_H^2$) is a signature of a strongly-interacting symmetry-breaking sector, since it would be absent if the Higgs boson were to be light. In contrast, the longitudinal gauge-boson interactions do not yet appear to be strong at these energies. The actual strong dynamics only shows up when the energy approaches the realm of the symmetry-breaking sector.

2.2. A systematic analysis of longitudinal gauge-boson scattering

We can now turn to a more general discussion of a strongly-interacting symmetry-breaking sector. To this end all other interactions (like weak and Yukawa interactions) are for the moment simply neglected. Two guiding principles are relevant for the discussion.

The first thing to note is that the SM Higgs sector has a larger symmetry than just the local $SU(2)_L \times U(1)_Y$. The Higgs Lagrangian

$$\mathcal{L}_H = (\partial_\mu \Phi)^\dagger (\partial^\mu \Phi) - \lambda \left(\Phi^\dagger \Phi - \frac{\mu^2}{2\lambda} \right)^2 , \tag{3}$$

with ϕ the complex Higgs doublet, is also symmetric under global $SU(2)_L \times SU(2)_R$ transformations. This corresponds to the symmetry under rotation of the four components. After the spontaneous symmetry breaking has taken place only the symmetry under rotation of the would-be Goldstone bosons is apparent. This (isospin) symmetry is usually referred to as the custodial $SU(2)_V$ symmetry. As far as strong interactions are concerned the same symmetry applies to the weak gauge bosons, being related to the would-be Goldstone bosons by the equivalence theorem. In this context both

(W^+, Z, W^-) and (ϕ^+, χ, ϕ^-) behave as isospin triplets ($I = 1$). In the SM this symmetry is (weakly) broken by hypercharge interactions. This leads to the relation $M_W = M_Z \cos\theta_w \neq M_Z$, with θ_w the weak mixing angle. The fact that the SU(2)$_V$ symmetry is not broken by the strong symmetry-breaking interactions is reflected by the observation that the so-called ρ parameter, representing the relative strength of neutral- and charged-current interactions at low energies, is close to unity: $\rho = 1 + \mathcal{O}(\%)$. In view of the strong experimental restrictions on this ρ parameter, any symmetry-breaking mechanism other than the one adopted in the SM should better obey the custodial SU(2)$_V$ symmetry.

The second guiding principle is provided by the LET corresponding to the spontaneous symmetry-breaking mechanism. The breaking of the (axial) symmetry, being of the order of the masses of the gauge bosons, is only a weak one when compared with the scale governing the strongly-interacting sector. In order to assess the implications of this observation, we first redefine the Higgs doublet by representing the would-be Goldstone bosons by a non-linear realization of the full SU(2)$_L \times$ SU(2)$_R$ symmetry group:

$$\Phi = \begin{pmatrix} \phi^+ \\ \frac{v+H+i\chi}{\sqrt{2}} \end{pmatrix} = \Sigma \begin{pmatrix} 0 \\ \frac{v+H'}{\sqrt{2}} \end{pmatrix} . \tag{4}$$

Here $\Sigma = \exp(i\omega^j \tau^j / v)$ transforms as $\Sigma \to \mathrm{U}_L \Sigma \mathrm{U}_R^\dagger$ under SU(2)$_L\times$SU(2)$_R$, with $\mathrm{U}_{L,R} \in$ SU(2) and τ^j ($j = 1, 2, 3$) the Pauli matrices. For energies well below M_H the heavy Higgs field can be integrated out, resulting in a non-renormalizable chiral Lagrangian:

$$\mathcal{L}_H = \frac{v^2}{4}\,\mathrm{Tr}\left([\partial_\mu \Sigma^\dagger][\partial^\mu \Sigma]\right) + \text{two terms with four derivatives} + \cdots \tag{5}$$

In this limit the would-be Goldstone fields ω^j are related to the original $\phi^\pm$ and χ fields according to

$$\omega^j = \phi^j \sum_{k=0}^{\infty} \frac{(2k)!}{2^{2k}(k!)^2(2k+1)} \left(\frac{\phi^l \phi^l}{v^2}\right)^k ,$$

$$\phi^\pm = \frac{\phi^1 \mp i\phi^2}{\sqrt{2}} \quad , \quad \chi = \phi^3 . \tag{6}$$

The chiral Lagrangian (5) represents the effective interactions between the would-be Goldstone bosons in the heavy-Higgs limit. It only involves derivative couplings, since $\Sigma^\dagger\Sigma = 1$. As a consequence there will be no strong scattering at low energies, *i.e.* the would-be Goldstone bosons decouple at low energies (up to gauge and Yukawa couplings). The first (kinetic) term in (5), with the lowest number of derivatives, is universal. Its coefficient is fixed by the electroweak symmetry-breaking scale v. The terms in the expansion with a larger number of derivatives are linked to the dynamics of the symmetry-breaking sector. They are suppressed by factors $E^2/(4\pi v)^2$, with $4\pi v$ the generic scale for the strong interactions (and resonances). In the case of the SM these terms will contain information on the Higgs sector. In general strongly-interacting scenarios the above chiral Lagrangian parametrizes the dynamics under the assumption of $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ symmetry, with the ω^j the would-be Goldstone bosons responsible for the generation of the gauge-boson masses.

What can we learn from longitudinal gauge-boson scattering, bearing in mind the above guiding principles for a general strongly-interacting symmetry-breaking sector? Applying the equivalence theorem, this is equivalent to an analysis of the generic process

$$\phi^i(p_i) + \phi^j(p_j) \to \phi^k(p_k) + \phi^l(p_l) \ , \tag{7}$$

involving identical massless spinless particles (as far as the strong interactions are concerned). Exploiting crossing symmetry and Bose symmetry for identical particles, the corresponding matrix element can be written as

$$\mathcal{M} = A(s,t,u)\,\delta_{ij}\delta_{kl} + A(t,s,u)\,\delta_{ik}\delta_{jl} + A(u,t,s)\,\delta_{il}\delta_{jk} \ , \tag{8}$$

with $A(s,t,u) = A(s,u,t)$. Here we introduced the standard Mandelstam variables $s = (p_i+p_j)^2$, $t = (p_i-p_k)^2 = -s(1-\cos\theta)/2$, and $u = (p_i-p_l)^2 = -s(1+\cos\theta)/2$, with $s+t+u=0$ and $\theta = \angle(\vec{p}_i,\vec{p}_k)$. Projection on the elastic (s-channel) isospin eigenstates yields:

$$\begin{aligned} I=0: &\quad \mathcal{M}_0 = 3A(s,t,u) + A(t,s,u) + A(u,t,s) \ , \\ I=1: &\quad \mathcal{M}_1 = A(t,s,u) - A(u,t,s) \ , \\ I=2: &\quad \mathcal{M}_2 = A(t,s,u) + A(u,t,s) \ . \end{aligned} \tag{9}$$

This can be rewritten in terms of charge eigenchannels:

$$\begin{aligned} \mathcal{M}(\phi^+\phi^- \to \phi^+\phi^-) &= A(s,t,u) + A(t,s,u) = \frac{1}{3}\mathcal{M}_0 + \frac{1}{2}\mathcal{M}_1 + \frac{1}{6}\mathcal{M}_2 \ , \\ \mathcal{M}(\phi^+\phi^- \to \chi\chi) &= A(s,t,u) = \frac{1}{3}\mathcal{M}_0 - \frac{1}{3}\mathcal{M}_2 \ , \\ \mathcal{M}(\chi\chi \to \chi\chi) &= A(s,t,u) + A(t,s,u) + A(u,t,s) = \frac{1}{3}\mathcal{M}_0 + \frac{2}{3}\mathcal{M}_2 \ , \\ \mathcal{M}(\phi^\pm\chi \to \phi^\pm\chi) &= A(t,s,u) = \frac{1}{2}\mathcal{M}_1 + \frac{1}{2}\mathcal{M}_2 \ , \\ \mathcal{M}(\phi^\pm\phi^\pm \to \phi^\pm\phi^\pm) &= A(t,s,u) + A(u,t,s) = \mathcal{M}_2 \ . \end{aligned} \tag{10}$$

For energies well below the scale of the strong symmetry-breaking interactions the LET predicts $A(s,t,u) = s/v^2 + \mathcal{O}(s^2/[16\pi^2 v^4])$. So, the amplitudes vanish at low energies (up to gauge and Yukawa couplings) with fixed slope at $s = 0$. The $\mathcal{O}(s^2/[16\pi^2 v^4])$ terms contain information on the symmetry-breaking dynamics.

The amplitudes for the elastic (s-channel) isospin eigenstates can be projected on partial waves:

$$\mathcal{M}_I = 32\pi \sum_{J=0}^{\infty} (2J+1)\, a_{IJ}(s)\, P_J(\cos\theta) \ , \tag{11}$$

with $P_J(\cos\theta)$ the Legendre polynomials and J the total angular momentum. Elastic unitarity reads $\mathrm{Im}(a_{IJ}) = |a_{IJ}|^2$ or $\mathrm{Im}(a_{IJ}^{-1}) = -1$ for all individual channels. This can be solved in terms of phase shifts: $a_{IJ} = \sin\delta_{IJ}\exp(i\delta_{IJ})$, spanning the unitarity circle. In the presence of inelastic channels, like $\phi\phi \to 4\phi$, the a_{IJ} are required to lie inside the unitarity circle, *i.e.* $\mathrm{Im}(a_{IJ}) > |a_{IJ}|^2$. These inelastic channels are suppressed in the energy expansion, since they only contribute at $\mathcal{O}(s^4/[4\pi v]^8)$. They are only relevant for energies above 2 TeV and are therefore neglected in the following. The LET predicts the lowest-order partial waves to be fixed:

$$a_{00} = \frac{s}{16\pi v^2} \ , \quad a_{11} = \frac{s}{96\pi v^2} \ , \quad a_{20} = -\frac{s}{32\pi v^2} \ . \tag{12}$$

Note that because of Bose symmetry $I+J$ should be even. The isoscalar ($I=J=0$) and isovector ($I=J=1$) partial waves are attractive, leaving open the possibility of finding a resonance at high energies in those channels. The partial wave for $I=2$ and $J=0$ is repulsive, excluding the presence of resonances.

In higher order in the energy expansion the partial waves are not universal anymore and two unknown parameters show up [9]: ($\varepsilon \downarrow 0$)

$$\begin{aligned} A(s,t,u) &= \frac{s}{v^2} + \frac{1}{16\pi^2 v^4}\left[\beta_1(\mu^2)\, s^2 + \beta_2(\mu^2)\, tu - \frac{1}{2}s^2\log\left(\frac{-s-i\varepsilon}{\mu^2}\right)\right. \\ &\left. - \frac{1}{6}t(t-u)\log\left(\frac{-t-i\varepsilon}{\mu^2}\right) - \frac{1}{6}u(u-t)\log\left(\frac{-u-i\varepsilon}{\mu^2}\right)\right] \ . \end{aligned} \tag{13}$$

Here the arbitrary parameter μ is merely introduced to make the arguments of the logarithms dimensionless. These logarithmic terms are a direct consequence of analyticity and elastic partial-wave unitarity for the a_{IJ}. In the language of chiral Lagrangians these logarithms are the result of chiral one-loop effects. For the renormalization of the one-loop effects the two terms in the chiral Lagrangian with four derivatives are required. This explains the occurrence of the unknown coefficients $\beta_{1,2}$, related to the dynamics of the model. In the same way the presence of the Higgs is required in the SM for having a renormalizable theory. The $\mathcal{O}(s^2/[16\pi^2 v^4])$ corrections to the partial-wave amplitudes a_{IJ} can be obtained from (9) and (13) by an appropriate projection of the matrix elements $\mathcal{M}_I$.

It turns out that the coefficients $\beta_{1,2}$ contribute with different signs for isoscalar and isovector resonances [10]. An isoscalar resonance gives rise to a positive contribution to a_{00} and a negative one to a_{11}. The reverse happens for an isovector resonance. As a result, the change in the slope of the partial waves provides crucial information on the strong dynamics, allowing a disentangling of the different models.

2.3. *Experimental sensitivity at the NLC*

At the NLC the longitudinal gauge-boson scattering processes show up in two distinct ways. The first one, displayed in Fig. 2a, involves the emission of massive gauge bosons (mainly W bosons) from the initial-state electron and positron. These gauge bosons subsequently interact with each other. This mechanism is called gauge-boson fusion. The advantage of this mechanism is the possibility to access all different channels (in isospin, angular momentum, and charge), especially if also the e^-e^- mode of the collider is used. By analyzing the invariant-mass distributions in the various channels, one should be able to differentiate between the various mass-generation models, provided the experimental resolution in the hadronic channels is sufficient for distinguishing between hadronic decays of W and Z bosons. The possibility to polarize the initial-state beams can be exploited to enhance the sensitivity, by increasing the number of W bosons emitted from the initial-state $e^\pm$. The main drawback of the fusion mechanism is the inefficient use of the collider energy, owing to the spectator leptons that carry away part

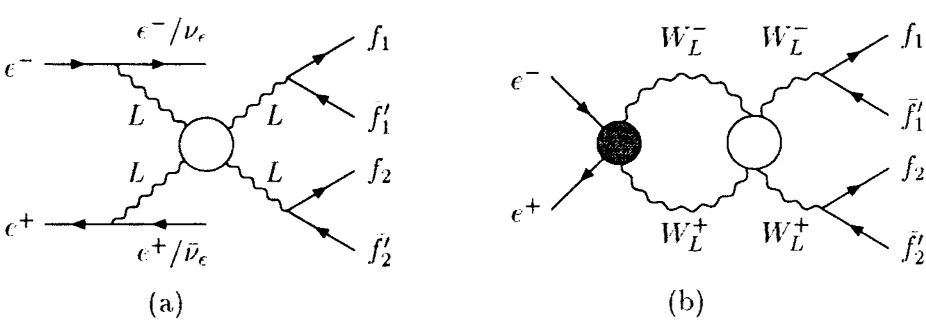

Fig. 2. The (a) fusion and (b) W-rescattering mechanisms for longitudinal gauge-boson scattering in e^+e^- collisions. The open circle connecting the longitudinal gauge-boson propagators represents the strong interactions.

of the energy. This explains why one has to resort to the invariant-mass distributions of the produced gauge bosons in order to investigate the dynamics of the symmetry-breaking sector; after all, the energy transmitted to the longitudinal gauge bosons is not fixed. As a result of the inefficient use of the collider energy, the signal cross-sections are relatively small until the actual resonances are formed.

The second and most promising way of studying longitudinal gauge-boson scattering at the NLC involves a detailed investigation of the process $e^+e^- \to W^+W^-$. The produced W bosons can trigger strong final-state interactions, called rescattering (see Fig. 2b). In view of angular-momentum conservation $J = 0$ is not allowed and only the $I = J = 1$ channel is accessible in this reaction. Making use of the elastic unitarity conditions for the strong final-state interactions in this isovector channel, the rescattering can be represented by a simple Muskhelishvili–Omnès form factor: ($\varepsilon \downarrow 0$)

$$\mathcal{M}(e^+e^- \to W^+W^-) = \mathcal{M}^{(0)} \exp\left[\frac{s}{\pi}\int_0^\infty ds'\, \frac{\delta_{11}(s')}{s'(s'-s-i\varepsilon)}\right] . \qquad (14)$$

Here $\delta_{11}(s)$ stands for the isovector phase shift and $\mathcal{M}^{(0)}$ indicates the lowest-order W-pair production amplitude without final-state interactions. When energy loss through initial-state photon radiation is kept under control, this rescattering process involves a relatively well-defined energy. The sensitivity to the interesting longitudinal gauge-boson modes can be enhanced by cutting away W bosons that are produced in the forward direction, which are predominantly transversely polarized. In addition, the angular distributions of the decay products of the W bosons can be exploited to increase the sensitivity to the longitudinal polarization states[2]. In order to have access to the Higgs-like $I = J = 0$ isoscalar channel at the NLC, one has to resort to the $\gamma\gamma$ collider mode with polarized photon beams. In this mode the isoscalar interactions can be investigated in the rescattering process $\gamma\gamma \to W^+W^-$ for energies comparable to the ones attainable with the e^+e^- mode. The overwhelming production of transverse W bosons, however, completely swamps the interesting rescattering phenomena. In this respect the process $\gamma\gamma \to ZZ$ looks more promising, but even there the transversely polarized Z bosons seriously hamper the study of a strongly-interacting isoscalar sector.

A recent study of the above processes [11], taking into account effects from initial-state radiation and beamstrahlung, has shown that a 500 GeV NLC with an integrated luminosity of 80 fb^{-1} will allow to exclude isovector

[2] It should be noted that the fermionic currents associated with these decays have (roughly) the same properties as the polarization vectors in (1), since a large majority of the decaying time-like gauge bosons is close to being on-shell.

resonances up to 2.5 TeV or discover such a resonance up to 1.5 TeV. A 1.5 TeV NLC with an integrated luminosity of 190 fb^{-1} should be able to compete with the Large Hadron Collider (LHC) in the isoscalar and non-resonant channels. In the isovector channel conclusive statements are expected: a strongly-interacting symmetry-breaking sector will be clearly distinguishable from the SM with a light Higgs, even if the associated isovector resonance has a very large mass. Even more, it will be possible to make statements concerning the mass of the isovector resonance. For instance, a 4 TeV resonance is expected to be distinguishable from an infinitely heavy 'resonance'. Based on these assessments it is safe to state that the NLC will be a prime machine for probing the symmetry-breaking sector, in particular if nature has chosen a strongly-interacting isovector scenario.

3. Physics beyond the Standard Model

The previous section has been exclusively dedicated to the mechanism of mass generation and its measurable effects at the NLC through longitudinal gauge-boson scattering. One may, however, ask oneself the question how any new physics (NP) beyond the SM will manifest itself. The most obvious signal would be the direct production of the particles associated with this NP sector. For this to happen the collider energy should be above the threshold for the production of these particles. If this is not the case the NP sector can only reveal itself indirectly, *i.e.* through deviations in the interactions between 'established' (SM) particles. These deviations are generally referred to as anomalous interactions. A parametrization of these anomalous interactions can be achieved by introducing the concept of effective Lagrangians.

3.1. The concept of effective Lagrangians and anomalous couplings

Let us assume that the energy scale associated with the NP sector (*e.g.* particle masses) is given by Λ_{NP} and that this energy scale largely exceeds the available collider energy. Then the NP effects will manifest themselves in two distinct (indirect) ways.

— *Exchange of heavy NP particles:* An example of an effective Lagrangian of this type is provided by the (pre-SM) Fermi contact interactions, describing the V$-$A structure of weak interactions at low energies. For example, the effective Lagrangian for the charged-current interactions between electrons, muons, and their neutrinos reads

$$\mathcal{L}_{CC}^{e,\mu} = -\frac{G_F}{\sqrt{2}}\, J_\lambda^\dagger J^\lambda \ , \qquad (15)$$

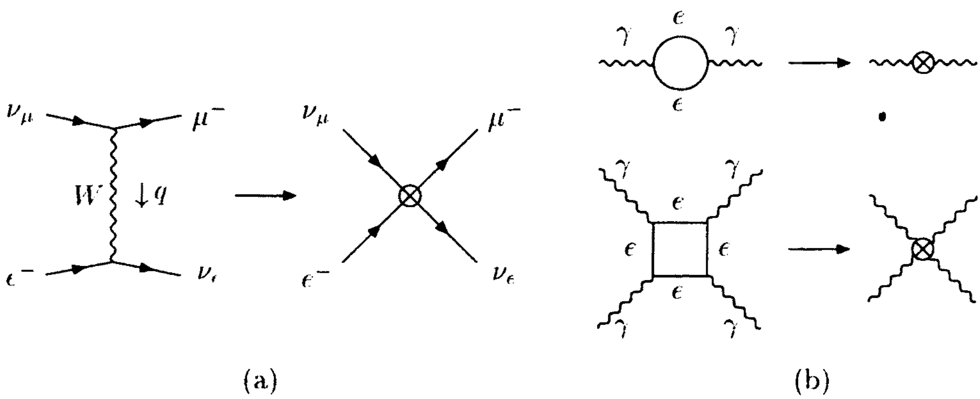

Fig. 3. Examples of (a) the exchange of heavy NP particles and (b) integrating out heavy NP particles in loops.

with

$$J^\lambda = \overline{\Psi}_{\nu_e}\gamma^\lambda(1-\gamma_5)\Psi_e + \overline{\Psi}_{\nu_\mu}\gamma^\lambda(1-\gamma_5)\Psi_\mu \ . \tag{16}$$

This effective Lagrangian is based on $U(1)$-invariant fermionic currents, as motivated by the (at that time) established theory of electromagnetic interactions. The Lagrangian $\mathcal{L}_{CC}^{e,\mu}$ is evidently not renormalizable. It is of dimension six, or in other words, the coupling constant G_F has dimension $(\mathrm{mass})^{-2}$. The cross-sections of reactions described by this effective Lagrangian seem to violate unitarity at high energies, *e.g.* $\sigma(\nu_\mu e^- \to \mu^-\nu_e) \sim G_F^2 s$. At high energies the underlying theory, *i.e.* the SM, will come to the rescue. To be more precise, the reaction is caused by the exchange of a spin-1 particle, called W boson (see Fig. 3a). The Fermi constant G_F will turn into a form factor $\sqrt{2}\,e^2/[8\sin^2\theta_w(M_W^2 - q^2)]$, ensuring that the cross-sections have a proper high-energy behaviour. From the low-energy limit $|q^2| \ll M_W^2$ one can deduce $G_F = \sqrt{2}\,e^2/[8\sin^2\theta_w M_W^2]$. At these energies the charged-current interactions appear weak, whereas for $|q^2| = \mathcal{O}(M_W^2)$ the full dynamics related to the SM W boson shows up and the electromagnetic and weak interactions become of comparable strength. Note also that in the above effective Lagrangian the concept of parity conservation, valid for the electromagnetic interactions, has been abandoned.

— *Integrating out heavy NP particles appearing in loops:* As an example we could go back to the pre-QED time. In this setting one could consider anomalous interactions between photons, caused by the interaction between these photons and unknown NP particles (called electrons). For energies $E_\gamma \ll m_e$ these low-energy interactions can be cast into an effective Lagrangian based on the free-photon field strength $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$:

$$\mathcal{L}_{\text{eff}} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \mathcal{L}_{\text{gauge fixing}} + \frac{\beta_1 e^2}{16\pi^2 m_\epsilon^2}\left[F^{\mu\nu}\Box F_{\mu\nu} + \frac{\epsilon_1}{m_\epsilon^2}F^{\mu\nu}\Box^2 F_{\mu\nu}\right]$$
$$+ \frac{e^4}{16\pi^2 m_\epsilon^4}\left[\beta_2(F^{\mu\nu}F_{\mu\nu})^2 + \beta_3(F_{\mu\nu}\varepsilon^{\mu\nu\rho\lambda}F_{\rho\lambda})^2\right] + \cdots \tag{17}$$

Again we end up with a Lagrangian that is non-renormalizable in finite order, *i.e.* at each order in the energy expansion the ultraviolet divergences can only be cancelled by introducing additional terms of higher order. The coefficients β_i and ϵ_i parametrize the deviations from the standard interactions and are ordered according to the expansion in E_γ/m_e or, equivalently, according to the dimension of the terms in the effective Lagrangian. At low energies the various terms in the effective Lagrangian seem to jeopardize unitarity. At higher energies, however, the dynamics of QED will show up, turning the coefficients into form factors and giving distinct predictions for the low-energy limits. These form factors will guarantee a proper high-energy behaviour and will be *related* according to the underlying theory (ensuring the renormalizability). Note that a factor $1/(16\pi^2)$ appears each time a heavy particle is integrated out in a loop. For instance, β_1 corresponds to the electron-loop contribution to the vacuum polarization, and $\beta_{2,3}$ correspond to the electron-loop contribution to light-by-light scattering (see Fig. 3b). The coefficient ϵ_1 corresponds to the energy expansion of the heavy-particle propagators appearing in the vacuum polarization and has accordingly no extra factor $1/(16\pi^2)$.

Bearing in mind the above examples, it is possible to construct a non-renormalizable effective Lagrangian describing physics beyond the SM. This effective Lagrangian takes the general form

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{NR}} \quad , \quad \mathcal{L}_{\text{NR}} = \sum_{n=5}^{\infty}\sum_i \frac{\alpha_i^{(n)}}{\Lambda_{NP}^{n-4}}O_i^{(n)} \; , \tag{18}$$

with n the dimension of the interaction, $\alpha_i^{(n)}$ the dimensionless anomalous couplings, and $O_i^{(n)}$ the operators describing the anomalous interactions between the 'established' particles. These operators respect the symmetry of the SM, *i.e.* they are invariant under the local $\text{SU}(2)_L \times \text{U}(1)_Y$ gauge transformations and depend only on covariant derivatives and field strengths. It should be noted that certain global symmetries that are present in the SM need not be maintained in the effective Lagrangian.

There are two scenarios for such an effective Lagrangian. In the first (linear) scenario the SM is completely established, including the presence of a light scalar Higgs boson. In that case the decoupling theorem [12] applies. This theorem states that if the heavy NP particles do not acquire their masses by means of the SM Higgs mechanism, then $\alpha_i^{(n)}$ does not depend on Λ_{NP}. Consequently the dimension-n operators are suppressed by factors $(Q/\Lambda_{NP})^{n-4}$, with Q denoting the masses of the SM particles or the energy of the collider (see the above examples). This introduces a natural hierarchy among the couplings that parametrize the NP effects at low energies. In the second (non-linear) scenario the Higgs is very heavy or absent altogether. In that case the effective Lagrangian will be based on $D_\mu \Sigma$ instead of $D_\mu \Phi$. As such it will resemble the chiral Lagrangian (5), except for the fact that the weak and NP couplings are not neglected with respect to the scale that governs the strongly-interacting symmetry-breaking sector. In this way a hierarchy emerges in powers of Q/Λ_{NP} and/or $E^2/(4\pi v)^2$.

3.2. Sensitivity to triple and quartic gauge-boson couplings at the NLC

Up to now only the gauge-boson–fermion couplings and the (bi-linear) couplings between two gauge bosons are tested with high precision at low-energy experiments like LEP1 and the SLC. The presence of NP effects in these couplings is excluded below the per-cent level. The natural next step would be to extend this to the non-abelian triple and quartic gauge-boson couplings. Such high-precision tests should be seen in the light of the afore-mentioned effective Lagrangian for NP effects, which will lead to specific contributions to the various couplings between the gauge bosons. The contributions to the triple gauge-boson couplings are strictly the result of integrated-out NP loop effects (leading to factors $1/16\pi^2$), whereas the contributions to the quartic gauge-boson couplings can also involve the exchange of NP particles.

A completely general investigation, involving the simultaneous effects of all possible gauge-boson couplings, is not recommendable in view of the expected statistics at the NLC. For instance, from angular-momentum conservation one can infer the existence of 14 independent couplings of the type $WW\gamma$ and WWZ [13][3]. In the actual data analysis one is going to adopt a more pragmatic attitude by only considering those interactions that are most likely to show up in the data. This opens the way to a large variety of theoretically and experimentally motivated prejudices, reducing the number of independent couplings. There are three main guiding principles. First of all, the ordering of the anomalous operators according to their dimension can be exploited by only taking into account the operators with

[3] In the derivation of this number the scalar parts of the off-shell gauge bosons were neglected. This is motivated by the fact that the gauge bosons are in general coupled to approximately massless fermionic currents.

the lowest dimension. For the gauge-boson interactions this boils down to restricting the analysis to dimension-six operators. Secondly, the charge of the W is fixed and the photonic interactions are too well-established to tamper with. Therefore C or CP violating $WW\gamma$ interactions can be discarded, and $U(1)_{em}$ gauge invariance should be preserved. Thirdly, the low-energy experiments strongly constrain certain couplings. Consequently, one should only consider operators that neither violate the custodial $SU(2)_V$ symmetry nor contribute to gauge-boson–fermion or bi-linear gauge-boson interactions. Combining all this one ends up with only three independent triple gauge-boson couplings. For more information the reader is referred to the literature [14].

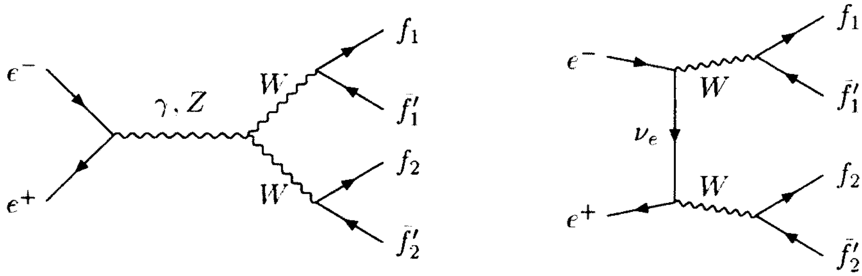

Fig. 4. The lowest-order 'W-pair' diagrams contributing to $e^+e^- \to f_1\bar{f}_1'f_2\bar{f}_2'$ in the SM. Here f_i' denotes the isospin partner of the fermion f_i. On the left: the s-channel diagrams involving the triple gauge-boson couplings. On the right: the t-channel ν_e-exchange diagram.

The main probe for anomalous triple gauge-boson couplings at the NLC will be the clean high-rate reaction $e^+e^- \to W^+W^-$ (see Fig. 4)[4]. Since the anomalous couplings are non-minimal, the delicate balance that is present between the SM diagrams is broken. For longitudinal gauge bosons this upsets the gauge cancellations at high energies (displayed in Fig. 5), leading to deviations that are enhanced by factors of order $\sqrt{s}/M_W$ for each longitudinal W boson. Based on this observation and the scaling properties of the anomalous interactions involving transverse W bosons, the sensitivity to the anomalous triple gauge-boson couplings is predicted to increase with the collider energy. As a rule of thumb the sensitivity attainable at e^+e^- colliders scales as $\sqrt{sL}$ [16], where L stands for the integrated luminosity. This automatically means that the NLC will do substantially better than LEP2 in view of the higher energy and luminosity. The sensitivity to the s-channel diagrams involving the triple gauge-boson couplings can be enhanced by cut-

[4] In [15] it has been shown that the sensitivity only marginally degrades when going from the full process $e^+e^- \to 4f$ to the 'W-pair' process $e^+e^- \to$ 'W^+W^-' $\to 4f$, which can be obtained by imposing tight invariant-mass cuts on the decay products.

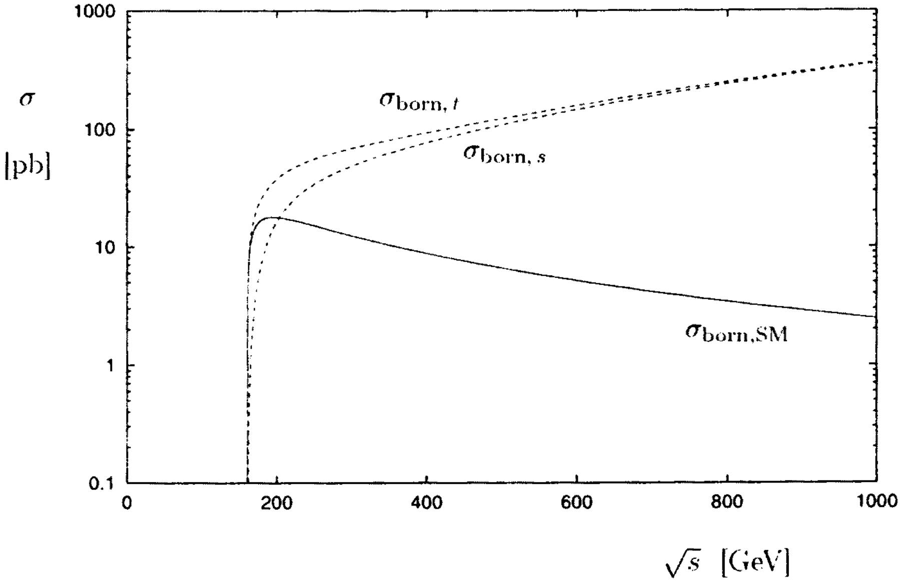

Fig. 5. Gauge cancellations in the lowest-order SM process $e^+e^- \to W^+W^-$. The dashed curves correspond to the Born cross-sections arising from the s-channel ($\sigma_{\rm born,s}$) and t-channel ($\sigma_{\rm born,t}$) diagrams alone. The solid curve corresponds to the complete Born cross-section ($\sigma_{\rm born,SM}$).

ting away the W bosons that are produced in the forward direction, since the forward direction is completely dominated by the contributions from the t-channel ν_e exchange. On top of that, in the clean environment of e^+e^- collisions one can use the angular distributions of the decay products of the W bosons as polarization analyzers. In this way the sensitivity to specific polarization states and hence to specific triple gauge-boson couplings can be enhanced. In order to fully exploit this opportunity it is important to identify the charge of at least one of the decay products of the decaying W bosons. At the NLC the possibility of having polarized initial-state beams allows to disentangle the effects from the $WW\gamma$ and WWZ couplings. In this context also the other collider modes come in handy, making a whole host of other reactions accessible.

Recent studies [15, 17] have shown that it will be feasible to probe the triple gauge-boson couplings at the permille level (at 500 GeV) or better (at 1.5 TeV). In this way the tests of the triple gauge-boson interactions would be promoted to the level of high-precision measurements. In order to test the quartic gauge-boson couplings one has to study gauge-boson fusion processes or processes that involve the production of three gauge-bosons (like $e^+e^- \to W^+W^-\gamma$, W^+W^-Z). In view of the reduced amount of available phase space and the suppression by the additional powers of the electroweak couplings, it will be difficult to beat the LHC at this job.

4. Accurate theoretical predictions in the SM

In order to successfully achieve the physics goals at the NLC, a very accurate knowledge of the SM predictions for the various observables is mandatory. This involves a proper understanding of radiative corrections as well as a proper treatment of finite-width effects.

4.1. The issue of gauge invariance

As has become clear from the previous sections, the W-boson physics studies at the NLC cover a large variety of processes with photons and/or fermions in the initial and final state. After all, the massive gauge bosons are unstable particles and can only be investigated through their decay products. If complete sets of graphs contributing to such a process are taken into account, the associated matrix elements are in principle gauge-invariant. However, the massive gauge bosons that appear as intermediate particles can give rise to poles $1/(k^2 - M^2)$ if they are treated as stable particles. This can be cured by introducing the finite decay width in one way or another, while at the same time preserving gauge independence and, through a proper high-energy behavior, unitarity. In field theory, such widths arise naturally from the imaginary parts of higher-order diagrams describing the gauge-boson self-energies, resummed to all orders. This procedure has been used with great success in the past: indeed, the Z resonance can be described to very high numerical accuracy. However, in doing a Dyson summation of self-energy graphs, we are singling out only a very limited subset of all the possible higher-order diagrams. It is therefore not surprising that one often ends up with a result that retains some gauge dependence.

Till recently two approaches for dealing with unstable gauge bosons were popular in the construction of lowest-order Monte Carlo generators. The first one involves the systematic replacement $1/(k^2 - M^2) \to 1/(k^2 - M^2 + iM\Gamma)$, also for $k^2 < 0$. Here Γ denotes the physical width of the gauge boson with mass M and momentum k. This scheme is called the 'fixed-width scheme'. As in general the resonant diagrams are not gauge-invariant by themselves, this substitution will destroy gauge invariance. Moreover, it has no physical motivation, since in perturbation theory the propagator for space-like momenta does not develop an imaginary part. Consequently, unitarity is violated in this scheme. To improve on the latter another approach can be adopted, involving the use of a running width $iM\Gamma(k^2)$ instead of the constant one $iM\Gamma$ ('running-width scheme'). This, however, still cannot cure the problem with gauge invariance.

At this point one might ask oneself the legitimate question whether the gauge-breaking terms are numerically relevant or not. After all, the gauge breaking is caused by the finite decay width and is, as such, in principle suppressed by powers of Γ/M. From LEP1 we know that gauge breaking can be negligible for all practical purposes. However, the presence of small scales can amplify the gauge-breaking terms. This is for instance the case for almost collinear space-like photons or longitudinal gauge bosons at high energies, involving scales of $\mathcal{O}(p_B^2/E_B^2)$ (with p_B the momentum of the involved gauge boson). In these situations the external current coupled to the photon or to the longitudinal gauge boson becomes approximately proportional to p_B. In other words, in these regimes sensible theoretical predictions are only possible if the amplitudes with external currents replaced by the corresponding gauge-boson momenta fulfill appropriate Ward identities.

In order to substantiate these statements, a truly gauge-invariant scheme is needed. It should be stressed, however, that any such scheme is arbitrary to a greater or lesser extent: since the Dyson summation must necessarily be taken to all orders of perturbation theory, and we are not able to compute the complete set of *all* Feynman diagrams to *all* orders, the various schemes differ even if they lead to formally gauge-invariant results. Bearing this in mind, we need some physical motivation for choosing a particular scheme. In this context two options can be mentioned, which fulfill the criteria of gauge invariance and physical motivation.

The first option is the so-called 'pole scheme' [18, 19, 20]. In this scheme one decomposes the complete amplitude according to the pole structure by expanding around the poles, *e.g.* $f(k^2)/(k^2 - M^2) = f(M^2)/(k^2 - M^2) +$ finite terms. As the physically observable residues of the poles are gauge-invariant, gauge invariance is not broken if the finite width is taken into account in the pole terms $\propto 1/(k^2 - M^2)$. It should be noted, however, that there exists some controversy in the literature [20, 21] about the 'correct' procedure for doing this and about the range of validity of the pole scheme, especially in the vicinity of thresholds.

The second option is based on the philosophy of trying to determine and include the minimal set of Feynman diagrams that is necessary for compensating the gauge violation caused by the self-energy graphs. This is obviously the theoretically most satisfying solution, but it may cause an increase in the complexity of the matrix elements and a consequent slowing down of the numerical calculations. For the gauge bosons we are guided by the observation that the lowest-order decay widths are exclusively given by the imaginary parts of the fermion loops in the one-loop self-energies. It is therefore natural to perform a Dyson summation of these fermionic one-loop self-energies and to include the other possible one-particle-irreducible fermionic one-loop corrections ('fermion-loop scheme') [22]. For the process $e^+e^- \to 4f$ (see Fig. 4), this amounts to adding the fermionic triple gauge-boson vertex corrections. The complete set of fermionic contributions forms a gauge-independent subset and obeys all Ward identities exactly, even with resummed propagators [23]. As mentioned above, the validity of the Ward identities guarantees a proper behavior of the cross-sections in the presence of collinear photons and at high energies in the presence of longitudinal gauge-boson modes. On top of that, within the fermion-loop scheme the appropriately renormalized matrix elements for the generic process $e^+e^- \to 4f$ can be formulated in terms of effective Born matrix elements, using the familiar language of running couplings [23].

A numerical comparison of the various schemes [22, 23] confirms the importance of not violating the Ward identities. For the process $e^+e^- \to e^-\bar{\nu}_e u\bar{d}$, a process that is particularly important for studying triple gauge-boson couplings, the impact of violating the electromagnetic $\mathrm{U}(1)_{\mathrm{em}}$ gauge invariance was demonstrated in [22]. Of the above-mentioned schemes only the running-width scheme violates $\mathrm{U}(1)_{\mathrm{em}}$ gauge invariance. The associated gauge-breaking terms are enhanced in a disastrous way by a factor of $\mathcal{O}(s/m_e^2)$, in view of the fact that the electron may emit a virtual (space-like) photon with p_γ^2 as small as m_e^2. A similar observation can be made at high energies when some of the intermediate gauge bosons become effectively longitudinal. There too the running-width scheme renders completely unreliable results [23]. In processes involving more intermediate gauge bosons, *e.g.* $e^+e^- \to 6f$ (see Fig. 2a), also the fixed-width scheme breaks down at high energies as a result of breaking $\mathrm{SU}(2)_L$ gauge invariance.

4.2. Radiative corrections in the SM

By employing the fermion-loop scheme all one-particle-irreducible fermionic one-loop corrections can be embedded in the tree-level matrix elements. This results in running couplings, propagator functions, vertex functions, etc. However, there is still the question about the bosonic corrections. Such corrections might mimic the presence of anomalous gauge-boson interactions if they are not taken into account properly[5]. A large part of the bosonic corrections, as *e.g.* the leading QED corrections, factorize and

[5] Note that the SM predictions need not have a 0.1% precision in order to study anomalous triple gauge-boson couplings at the 10^{-3} level. These anomalous couplings are non-minimal and therefore lead to enhanced effects.

can be treated by means of a convolution, using the fermion-loop-improved cross-sections in the integration kernels (see *e.g.* appendix A of [24]). This allows the inclusion of higher-order QED corrections and soft-photon exponentiation. In this way various important effects can be covered. In this respect especially the emission of hard photons from the initial state is noteworthy [25, 26]. The associated hard-photon boost effects will lead to a redistribution of phase space, which affects the angular distributions. Also the polarization of the produced gauge bosons is affected by the presence of such boosts. The best-suited observables for probing the NP sector normally involve small cross-sections and are therefore extremely sensitive to redistribution effects. For instance, observables related to longitudinal gauge bosons will receive large corrections from the tranverse modes. In order to reduce the effects from the hard photons to a minimum, appropriate cuts on the total final-state energy and momentum have to be imposed.

Only taking into account the (leading) factorizing bosonic corrections is not sufficient. The remaining bosonic corrections can be large, especially at high energies where logarithmic corrections $\propto \log^2(M^2_{W,Z}/s)$ emerge [26, 27]. In order to include the remaining bosonic corrections one might attempt to extend the fermion-loop scheme. In the context of the background-field method a Dyson summation of bosonic self-energies can be performed without violating the Ward identities [28]. However, the resulting matrix elements depend on the quantum gauge parameter at the loop level that is not completely taken into account. As mentioned before, the perturbation series has to be truncated; in that sense the dependence on the quantum gauge parameter could be viewed as a parametrization of the associated ambiguity.

As a more appealing strategy one might adopt a hybrid scheme, adding the remaining bosonic loop corrections by means of the pole scheme. This is gauge-invariant and contains the well-known bosonic corrections for the production of on-shell gauge bosons (in particular W-boson pairs [29, 30]). Moreover, if the quality of the pole scheme were to degrade in certain regions of phase-space, the associated error is reduced by factors of α/π. It should be noted that the application of the pole scheme to photonic corrections requires some special care, because in that case terms proportional to $(k^2 - M^2)^{-1}\log(k^2 - M^2)$ complicate the pole expansion [20, 26].

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