Anomalous high-energy behaviour in boson fusion

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Abstract

We discuss the magnitude of the off-shell correction terms to the equivalent-boson approximation used in SSC phenomenology, with emphasis on Higgs production. Although formally subleading, these terms are NOT small and give rise to cross sections that increase unacceptably at high invariant mass of the boson pair. We conclude that only a much more complete set of Feynman diagrams than those used in the equivalent-boson approximation can yield reliable predictions for the production of a heavy Higgs boson.

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This letter deals with the phenomenology of electroweak interactions at an energy scale of the order of a TeV, in particular the production of a heavy Higgs boson at the next generation of hadron colliders [1]. It is clear that if this process is to provide a probe of the standard model, then the predictions for its cross section must be as accurate as possible. So far, calculations of Higgs production have relied on one or more approximations. Cahn and Dawson [2], [3] studied the production of an on-shell, non-decaying Higgs boson via the mechanism of WW fusion [4]. They made the approximation that the W's are produced as on-shell 'partons' from quark scattering over negligible angle. If the Higgs is light (and therefore long-lived) the use of this single Feynman graph is more or less justified. However, for low values of the Higgs mass the WW fusion is by no means the most important mechanism, gluon fusion having a comparable rate [2]. A slightly more sophisticated analysis [5], [6] uses the graph of fig.1^a, where the Higgs decay into, say, a pair of Z's is taken into account. In doing so, however, one has to be quite careful: as was pointed out by Duncan, Kane and Repko [7] such a diagram is not gauge-invariant, and has bad high-energy behaviour; in fact it forms part of a larger, gauge-invariant group of diagrams. The authors of ref [7] studied the behaviour of such a set of diagrams, namely those describing $W^+W^- \rightarrow W^+W^-$; we briefly recapitulate their results. In the limit that s , the invariant mass of the WW pair, goes to infinity, the total amplitude must go to a constant as required by unitarity. For longitudinally polarized W's, however, the diagrams containing the three- and four-boson vertices grow like s^4/m^4 in this

limit (assuming the Feynman gauge), where m stands for the W mass; their sum exhibits cancellations and grows as s/m^2 ; only upon including the two Higgs graphs does the last power of s cancel and the correct high-energy behaviour obtains. As we have checked, identical behaviour is observed for the slightly simpler and phenomenologically more attractive [5], [6], [8] case of $WW \rightarrow ZZ$. The relevant set of diagrams is depicted in fig. 1^{a-d}. In view of the large width expected for a heavy Higgs, which weakens the 'resonant' behaviour of s -channel Higgs exchange, it is clear that all these graphs have to be taken into account in a phenomenological study with any claim to accuracy [9].

All studies so far have used the approximation that the incoming W 's or Z 's are on-shell particles. This is usually justified by the argument that the W 's are quite light (about .1 TeV) compared to the typical scale of the scattering process (up to a few TeV); and the propagators of the virtual bosons are indeed maximal at zero scattering angle, making the expected typical q^2 in these propagators also small. This argument is known to work quite well in two-photon physics in e^+e^- collisions; the massless virtual photons indeed tend to have a q^2 very close to 0, and the equivalent-photon approximation is known to be good to about the 10% level; a similar accuracy for the equivalent-boson approximation was reported in [10].

We would like to point out, however, that it may be dangerous to blindly carry over the experiences from two-photon physics to two-boson physics. A rough, heuristic argument which illustrates this is as follows. In $\gamma\gamma$ scattering a typical cross section involves an integral over the q^2 values of the virtual photons:

$$\sigma \sim \int \frac{dq_1^2}{q_1^2} \frac{dq_2^2}{q_2^2} \sigma_{sub}(s, q_1^2, q_2^2) \quad , \quad (1)$$

where σ_{sub} is the subprocess cross section for $\gamma\gamma$ scattering into the final state. This we take to be a two-body final state, say, an electron pair. As stated above, at high s σ_{sub} should go to a constant for 2→2 scattering. Assuming reasonably that $q_1^2, q_2^2 \ll s$ for high s^1 we can write

$$\sigma_{\text{sub}}(s, q_1^2, q_2^2) = C_1 + C_2 \frac{q_1^2}{s} + C_3 \frac{q_2^2}{s} + \dots, \quad (2)$$

since s is essentially the only scale in a QED process like $\gamma\gamma \rightarrow e^+e^-$. The coefficients C_1, C_2 and C_3 are numbers of more or less the same magnitude, that go to constants as $s \rightarrow \infty$. C_1 is the cross section for on-shell photons. Performing the integral (1) with boundaries $Q_0^2 \approx 0$ and $|Q_1^2| \ll s$, respectively, we obtain

$$\sigma \sim C_1 \left[\ln \frac{Q_1^2}{Q_0^2} \right]^2 + \frac{Q_1^2}{s} (C_2 + C_3) \ln \frac{Q_1^2}{Q_0^2} + \dots \quad (3)$$

We see that the term with C_1 gives the leading-log behaviour, and the coefficients of the subleading terms are of the same order of magnitude as the one of the leading term.

The analogous expression for the WW scattering case is

$$\sigma \sim \int \frac{dq_1^2}{(q_1^2 - m^2)} \frac{dq_2^2}{(q_2^2 - m^2)} \sigma_{\text{sub}}(s, q_1^2, q_2^2) \quad (4)$$

¹ This is not automatically satisfied since s is the invariant mass of the subprocess only; the physical limit on q^2 is rather s_0 , the total c.m. energy of the colliding e^+e^- pair.

In writing out σ_{sub} we now have to be more careful. Since it is not a priori clear that σ_{sub} has good high-energy behaviour for OFF-SHELL bosons, other terms may arise. In particular, since now not only s , but also the boson mass m is a scale in the problem, we have a dimensionless number, s/m^2 , that may enter in the coefficients C_2 , C_3 , ..., leading to

$$\sigma_{\text{sub}}(s, q_1^2, q_2^2) = C_1 + \left(\frac{s}{m^2}\right)^{n_2} C_2 \frac{q_1^2 - m^2}{s} + \left(\frac{s}{m^2}\right)^{n_3} C_3 \frac{q_2^2 - m^2}{s} + \dots \quad (5)$$

with some powers $n_2, n_3, \dots$; this leads to an integral like

$$\sigma \sim C_1 \left[\ln \left| \frac{Q_1^2}{m^2} \right| \right]^2 + \frac{Q_1^2}{s} \left[\left(\frac{s}{m^2}\right)^{n_2} C_2 + \left(\frac{s}{m^2}\right)^{n_3} C_3 \right] \ln \left| \frac{Q_1^2}{m^2} \right| + \dots \quad (6)$$

That powers n_i different from 0 may arise is indicated by the above strong and delicate cancellations between the various graphs for boson-boson scattering; if such cancellations are not complete in the off-shell case, then terms like $s(q_1^2 - m^2)/m^4$ can arise. If so, the behaviour of the cross section as $s \rightarrow \infty$ may be irreparably wrong, even with suppression of the off-shell terms by the boson propagators. As we shall argue later, this is indeed the case.

The above argument is by no means meant to be rigorous. With it we only want to indicate that a potential source of problems is uncovered as soon as we take off-shell effects into account; but, in order to have an accurate phenomenological prediction, they have to be considered:

In order to check the above arguments we have studied the behaviour of the boson-boson amplitudes in the case where the incoming bosons are longitudinally polarized. This is known to be the polarization amplitude with the worst high-energy behaviour [7]. Our approach is the following. Taking a process where two bosons collide in their centre-of-mass frame with longitudinal polarization, we evaluate the leading high-energy behaviour of the amplitude. This we do for a process which is known to be well-behaved at high energies, like $\gamma\gamma \rightarrow e^+e^-$, and then again for processes like $W^+W^- \rightarrow Z^0Z^0$. Comparing these results we can judge whether bad high-energy behaviour is to be expected, or not. We do this in three cases: both bosons off-shell, one on- and one off-shell, and both on-shell. In the last case the high-energy behaviour is known to be reasonable, which provides a check on our procedure.

Following this approach, we avoid having to worry about the luminosity functions for polarized W's, Z's and γ 's, and the influence of the propagators; these are essentially the same for all types of bosons. An additional advantage is that for the 2 $\rightarrow$ 2 subprocess the amplitude is relatively simple and can be studied analytically, using computer algebra. Of course, ultimately the bad high-energy behaviour must be exhibited in an explicit numerical calculation; this we shall do below.

As an example, we discuss the amplitude for

$$W^+(p_1) W^-(p_2) \rightarrow Z^0(p_3) Z^0(p_4) \quad , \quad (7)$$

where all the bosons have longitudinal polarization, as indicated above. The Z-boson pair is assumed to be on-shell: $p_3^2 = p_4^2 = m_Z^2$. The W pair is off-shell:

$$p_{1,2}^2 = m^2 + 4E \Delta_{1,2} \quad , \quad (8)$$

where $2E$ is the centre-of mass energy of the scattering process. The variables Δ_1, Δ_2 parametrize the off-shellness of the bosons; by letting them go to 0 we put the bosons on-shell. In the physical case, however, they have an upper limit $-m^2/(4E)$ since p_1^2 and p_2^2 must be negative. The momenta in the centre-of-mass frame are as follows:

$$\begin{aligned} p_1^\mu &= (E + \Delta_1 - \Delta_2, \quad p_w \vec{e}_w) , \\ p_2^\mu &= (E - \Delta_1 + \Delta_2, \quad -p_w \vec{e}_w) , \\ p_3^\mu &= (E, \quad p_z \vec{e}_z) , \\ p_4^\mu &= (E, \quad -p_z \vec{e}_z) ; \\ p_w &= [(E + \Delta_1 - \Delta_2)^2 - 4E\Delta_1 - m^2]^{1/2} , \quad p_z = [E^2 - m_z^2]^{1/2} , \end{aligned} \quad (9)$$

where $\vec{e}_w$ and $\vec{e}_z$ are unit three-vectors. The scattering angle θ is given by $\cos\theta = \vec{e}_w \cdot \vec{e}_z$. The polarization vectors, denoted by ϵ_i^μ , are

$$\begin{aligned} \epsilon_1^\mu &= (p_w, \quad [E + \Delta_1 - \Delta_2] \vec{e}_w) , \\ \epsilon_2^\mu &= (p_w, \quad -[E - \Delta_1 + \Delta_2] \vec{e}_w) , \\ \epsilon_3^\mu &= (p_z, \quad E \vec{e}_z) / m_z , \\ \epsilon_4^\mu &= (p_z, \quad -E \vec{e}_z) / m_z . \end{aligned} \quad (10)$$

For the Z's, these are what one would expect for a longitudinal polarization, being normalized to -1; but for the W's we have instead $\epsilon_i \cdot \epsilon_i = -p_i^2 > 0$, instead of -1. Using this normalization, however, we automatically in-

clude the fact that the luminosity function² for longitudinally polarized bosons is proportional to p_i^2 , where p_i^μ is the virtual boson momentum. As a consequence the dimension of the 2→2 scattering amplitude will now be (mass)² instead of (mass)⁰.

The amplitude for the process (7) is given by

$$\begin{aligned}
 M = -ig^2 \Big\{ & \cos^2\theta_w \Big\{ \left[2(p_i \cdot \epsilon_3) \epsilon_i^\alpha + 2(p_3 \cdot \epsilon_i) \epsilon_3^\alpha - (\epsilon_i \cdot \epsilon_3)(p_i + p_3)^\alpha \right] \times \\
 & \left[2(p_2 \cdot \epsilon_4) \epsilon_2^\beta + 2(p_4 \cdot \epsilon_2) \epsilon_4^\beta - (\epsilon_2 \cdot \epsilon_4)(p_2 + p_4)^\beta \right] \times \\
 & \left[g_{\alpha\beta} - (p_i - p_3)_\alpha (p_i - p_3)_\beta / m^2 \right] / [(p_i - p_3)^2 - m^2] \\
 & + (p_3, \epsilon_3 \leftrightarrow p_4, \epsilon_4) \Big\} \\
 & + \cos^2\theta_w \left[2(\epsilon_i \epsilon_2)(\epsilon_3 \epsilon_4) - (\epsilon_i \epsilon_3)(\epsilon_2 \epsilon_4) - (\epsilon_i \epsilon_4)(\epsilon_2 \epsilon_3) \right] \\
 & + \frac{m_Z^2}{2} (\epsilon_i \epsilon_2)(\epsilon_3 \epsilon_4) / [(p_3 + p_4)^2 - m_H^2] \Big\}, \quad (11)
 \end{aligned}$$

where g is the weak coupling constant $e/\sin\theta_w$, and θ_w is the weak mixing angle. m_H is the Higgs mass, and for the Z mass the relation $m_Z = m/\cos\theta_w$ holds. We have evaluated this amplitude using REDUCE [11]. Since the resulting expression contains some 600 terms we do not present it here, but its leading high-energy behaviour can easily be determined. If both Δ_1 and Δ_2 are nonzero, the amplitude behaves as $E^4 \Delta_1 \Delta_2 / m^4$, where we

² Due to the different normalization of the polarization vectors, our definition of the luminosity function is not identical to, say, the one employed in [7].

have dropped all coupling constants, powers of $\sin^2\theta_W$ and angular dependences (incidentally, in this limit the amplitude happens to be independent of θ). If we let Δ_2 go to zero while retaining a nonzero Δ_1 , the leading behaviour is $E^3\Delta_1/m^2$, and if both Δ 's are zero the amplitude reduces to (m_H^2/m^2) , with which we indicate a linear combination of m_H^2 and m_W^2 . In this latter case, the actual longitudinal polarization vectors for on-shell W 's differ from those given in eq.(10) by a factor of $1/m$; the on-shell amplitude is therefore seen to be $(m_H^2/m^2; 1)$, i.e. a dimensionless constant, as required.

We have examined various processes in the above manner. We give the leading behaviour in each of the three possible cases in the table. From this table we can draw the following conclusions. The process $ZZ \rightarrow ee$ is wellbehaved; in fact, setting m_Z to zero we obtain the same result, now valid for $\gamma\gamma \rightarrow ee$. The process $WW \rightarrow ee$ behaves slightly less well, having an amplitude that grows with energy. However, if we realize that $p^2 = m^2 + 4E\Delta$ we see that for fixed p^2 the amplitude is actually constant. Therefore, we do not expect bad high-energy behaviour for this process. The process $ZZ \rightarrow ZZ$ is quite wellbehaved; the amplitude remains reasonable even if both the incoming Z 's are far off-shell. This could also be guessed from the fact that for this process, ONLY Higgs exchange contributes (in both s -, t - and u -channel); since these graphs grow only like s/m^2 (as discussed above) no large cancellations are present, and hence there is not much opportunity to spoil them. The processes $WW \rightarrow ZZ$, $ZZ \rightarrow WW$ and $WW \rightarrow WW$, however, exhibit dramatically bad high-energy behaviour as compared to the 'good' processes. Even when we require the virtual bosons to have zero q^2 (i.e. $\Delta_1, \Delta_2 = -m^2/4E$), the amplitude is still rising as E^2 , and the cross section as E^4 . It has already been argued [12] that the high-energy be-

Table			
High-energy behaviour of the scattering amplitude			
Each incoming boson is either off-shell or on-shell			
process	off-off	on-off	on-on
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ZZ→ee	$\Delta_1 \Delta_2$	$m^2 \Delta_1 / E$	m^4 / E^2
WW→ee	$E(\Delta_1 + \Delta_2)$	$E \Delta_1$	m^2
ZZ→ZZ	$(m_H^2; \Delta_1 \Delta_2)$	m_H^2	m_H^2
ZZ→WW	$\Delta_1 \Delta_2 E^4 / m^4$	$\Delta_1 E^3 / m^3$	$(m_H^2; m^2)$
WW→ZZ			
WW→WW			

haviour would probably be under control if the bosons are not too far off shell. As can be seen from the table, a restriction like $p_{1,2}^2 = m^2 + \mathcal{O}(m^3/E)$ would indeed lead to an amplitude that goes as m^2 for $E \rightarrow \infty$; unfortunately, this q^2 range happens to be outside the accessible phase space. We infer from these results that for these three processes the high- s behaviour is anomalous. What is more, the bad behaviour persists even for decent values of Δ_1 and Δ_2 . We therefore expect that there is no practicable cut on the event topology which makes these cross-sections wellbehaved.

The source of the problems is easily seen to be the presence of the nonabelian vertices between 3 or 4 bosons. It is these vertices that are responsible for the strong rise of these diagrams with s , and make subtle cancellations between them necessary (as can be seen from the case $ZZ \rightarrow ZZ$, the Higgs exchange graphs do not cause any problem). The reasonable agreement between exact and approximate results reported in ref [10] can be traced to the fact that there the Higgs was produced as a stable particle in the final state: consequently, no 3-boson vertices occurred in that calculation. We conjecture therefore that similar or even worse behaviour will be observed for a process like $WW \rightarrow WWZ$.

The above reasoning may be convincing, but of course it has to be supported by an exact numerical calculation, taking also the other polarization states of the bosons and the various corresponding luminosity functions into account. To this end we have written a Monte Carlo program which integrates the complete Feynman diagrams for $WW \rightarrow ZZ$ and $ZZ \rightarrow ZZ$, including the (anti)quark lines that emit the incoming bosons, and the fermionic decays of the outgoing bosons. It will be described in greater detail elsewhere [13]: here we just mention its most important features. The incoming and outgoing bosons are described by Breit-Wigner resonances; in the decays of the produced Z's the correct decay correlations are taken into account. This is facilitated by the use of spinor techniques [14]. Although this is not very important to our present discussion it enables one to impose realistic cuts and influences the detailed phenomenology [6] [8]. The fact that we describe the boson emission from the incoming quark lines exactly has the advantage that one can in principle make detailed studies of, for example, the rapidity and transverse momentum distributions

of the scattered partons [6]. It also implies that the virtual bosons are described by the correct density matrices: no approximation as for equivalent bosons is necessary, and both the longitudinal and transverse degrees of freedom contribute coherently. Finally, the structure functions are those of ref [15], which are especially suited to the high energies of SSC.³ Throughout our discussion we use a total SSC energy of 40 TeV. The Higgs mass is arbitrarily taken to be 400 GeV, hence its width is 25 GeV in the Standard Model.

We start by presenting the results for the 'wellbehaved' process $ZZ \rightarrow ZZ$. In fig 2^a the distribution of the invariant mass m_{ZZ} of the Z pair is given. The Higgs resonance stands out as a clear peak at $m_{ZZ}=m_H$. At higher values of m_{ZZ} the cross section falls uniformly. The distribution in s' , the total c.m. energy of the Z pair and the scattered (anti)quarks, shows a peak around 1.5 TeV after which it falls in a similar manner (fig.2^b). The $\sqrt{(-q^2)}$ distribution (fig 2^c) of the incoming Z's peaks around 100 GeV which is indeed of the same order as m_Z . In all, the cross section is seen to behave as one would reasonably expect. We shall not comment on any further phenomenological details here but reserve these for a future publication [13].

We now turn to the process $WW \rightarrow ZZ$. In fig.3 we give the results of the integral over the complete phase space, without any cuts. As predicted by the above qualitative considerations, the distribution of m_{ZZ} (fig.3^a) is

³ We have taken a scale of 200 GeV. This is somewhat arbitrary; however taking a different scale does not change our conclusions.

seen to be anomalous, rising throughout the Higgs resonance region: no Higgs resonance is evident. The actual maximum is reached between 3 and 4 TeV, with an appreciable tail extending as far as 16 TeV. Similar behaviour is observed for the s' distribution (fig.3^b) which peaks at around 14 TeV with a tail up to 28 TeV. The $\sqrt{(-q^2)}$ distribution (fig.3^c) is quite interesting: after an initial peak around 100 GeV it first falls by about a factor 3 at 250 GeV, and then starts rising again, to reach several times the value of the first peak around 3-4 TeV. Clearly this behaviour is anomalous; there is nothing in this set of diagrams to make the various distributions peak at these high values. In fact the distributions only decrease again because of the suppression by the structure functions. This is indicated by the fact that if we increase the overall energy from 40 TeV to 80 TeV the peaks are also displaced by a factor two. It might be thought that this is related to the fact that the $\sqrt{(-q^2)}$ distribution begins to rise above 250 GeV. In that case one might expect good behaviour by making a maximum- q^2 cut around $-(200 \text{ GeV})^2$. The effect of this on the m_{ZZ} distribution is shown in fig 4. Because of the rejection of a large part of the cross section, the Higgs resonance can now just be distinguished. But the rising behaviour of m_{ZZ} is the same as before. In particular, the maximum of its distribution remains at the same value (the same holds for the s' distribution).

We have shown that both (qualitative) analytic and exact numerical investigation indicates that the high-energy behaviour of these bosonic processes is anomalous. But, clearly, unitarity must be satisfied, and the cross section must remain finite at all energies. We now proceed to solve this dilemma.

The source of the anomalous behaviour was seen to be the fact that the incoming bosons, being emitted in the t -channel with negative q^2 , are always off-shell by at least as much as their own mass. But, being off-shell, they cannot be the asymptotic states in the problem: this is in fact taken into account by representing the bosons as Breit-Wigner resonances emitted by the (anti)quark lines. But then, we formally have to consider our process as not $WW \rightarrow ZZ$ but rather as $q_1 q_2 \rightarrow q'_1 q'_2 ZZ$ (where the q 's stand for the (anti)quarks), and consider ALL the Feynman diagrams describing this process, ALSO the ones that cannot easily be interpreted as describing boson fusion. In fig.5 we indicate this much larger set of diagrams. The blob represents the $WW \rightarrow ZZ$ scattering given in fig.1; the other diagrams are added to these ones coherently. If q_1 is an up-quark, and q_2 a down-quark, there are 28 of these extra diagrams. We hold that only upon including these extra graphs can the desired high-energy behaviour be obtained. We have checked that this is true if one of the incoming W 's is on-shell. In the process $q W \rightarrow q' W Z Z$ the sum of the 4 diagrams in which both Z 's are emitted by the W line diverges at high energy; the remaining 10 diagrams, with either one or both of the Z 's coming from the quark line, give a contribution that precisely cancels this unwanted high-energy behaviour. This problem will be discussed more fully elsewhere [13].

We conclude this paper by summarizing the results. We have argued that in some processes mediated by the fusion of two virtual vector bosons (W 's or Z 's), in particular the production and decay of a heavy Higgs boson, the cross section may have bad high-energy behaviour as a consequence of the presence of many Yang-Mills couplings. This is confirmed

by both analytic and numerical studies. We furthermore argue that this problem can only be solved by taking into account a far larger set of Feynman diagrams, many of which do not contain the two virtual bosons present in the original graphs. Unfortunately, this shows that more work is needed before the boson fusion picture can be considered a well-justified approximation.

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1. Figure captions

fig.1

Feynman diagrams for the process $W^+W^- \rightarrow Z^0Z^0$

fig.2

Distributions for the process $Z^0Z^0 \rightarrow Z^0Z^0$: (a) the invariant mass m_{ZZ} of the Z pair; (b) the invariant mass s' of the original (anti)quark pair; (c) the momentum transfer $\sqrt{(-q^2)}$ of the radiated virtual bosons.

fig.3

The same as fig.2 but now for the process $W^+W^- \rightarrow Z^0Z^0$.

fig.4

The same as fig.3^a but now with a cut on $\sqrt{(-q^2)}$ of 200 GeV.

fig.5

Feynman diagrams for the process $q_1q_2 \rightarrow q'_1q'_2Z^0Z^0$. The blob in diagram (a) contains the diagrams of fig.1. There are 8 diagrams of type (b), and 20 diagrams of type (c).

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5