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Effective Lagrangians and $W_L W_L$ Scattering

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ABSTRACT

In this talk we discuss how to use effective chiral Lagrangians to investigate in some detail Chanowitz's "no lose corollary". That is, we will assume that electroweak symmetry breaking occurs through an unknown strong interaction that, however, produces no resonances in the energy region that will be probed by a *pp* supercollider like the SSC. We find that if an enhancement in the yield of $V_L V_L$ pairs is observed, it will be very difficult to relate it to an underlying theory. We also find that in a "worse case scenario" this enhancement might not be sufficiently large for detection at the SSC.

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1. Introduction

One of the major reasons for building high energy hadron colliders is to study the mechanism of electroweak symmetry breaking. Since it is this mechanism that is responsible for giving the W and Z gauge bosons their masses, (and therefore their longitudinal degrees of freedom) one expects the interactions of longitudinally polarized vector bosons to be particularly sensitive to it.

A subject that has received much attention, is the production and observation of new particles associated with the breaking of electroweak symmetry. The case that has been studied most is, of course, the standard model with an elementary Higgs boson. [1] If a Higgs boson is not found in the near future, the question arises of how large can the Higgs boson mass be. It has been argued that beyond a certain value of the Higgs mass, typically around 800 GeV, the symmetry breaking sector of the minimal standard model becomes strongly interacting. [2] One could still have a resonance with the quantum numbers of the Higgs, although its properties would not be simply related to the parameters in the model. Such a strongly interacting symmetry breaking sector need not be at all like the minimal standard model, and many alternatives have been suggested. In that case one could also have a richer spectrum of resonances in the few TeV region as is expected for technicolor models. [3]

Ideally, the study of WW scattering will reveal the presence of any new resonance associated with the physics responsible for the symmetry breaking (as a Higgs boson, a techni-rho, etc.). It has been claimed that even if there are no such resonances at sufficiently low energies to be detected at the SSC, one would observe a significant enhancement in the yield of WW pairs, signaling a strongly interacting symmetry breaking sector. This is the so called "no lose corollary". [4] It is with this latter scenario that we will concern ourselves. [5]

The interest is then to study WW scattering in a "low energy region", below threshold for production of any new resonance. A particularly useful framework for such studies is that of effective Lagrangians and chiral perturbation theory, which allows one to study general features in a model independent way. Within that approach, one parameterizes the dynamics of the new physics in terms of a few unknown coupling constants. [6]

The questions that we will address are two. First we will assume that there are indeed new resonances associated with the breakdown of electroweak symmetry, but that they are beyond the reach of the SSC. The physics that originates them would, however, show up at "low energies" in the form of somewhat enhanced production of WW pairs. Qualitatively we would like to know whether it is possible to learn about unseen resonances by studying the accessible energy region. [7] The second question that we will address is the generality of the "no lose" corollary. That is, we will want to know if for an arbitrary theory with strongly interacting vector bosons there is indeed an enhanced yield of WW pairs that could be detected at the SSC. [8]

We will use the equivalence theorem [9] to extract from the full amplitude only those terms that are of "enhanced electroweak strength". This means that we will calculate our amplitudes replacing all W and Z gauge bosons by their corresponding would-be Goldstone bosons in Landau gauge. By doing this we will obtain amplitudes that are correct up to terms of order $\mathcal{O}(M_W/E)$. Our results for WW scattering will thus be valid only for $s \gg M_W^2$. One could relax this condition by calculating exactly (instead of using the equivalence theorem). However, it is at high energies that one expects the "symmetry breaking" effects to become important.

The processes that will occur at the SSC are of course very complicated. One has two protons colliding that yield complicated final states with leptons and jets. For our discussion we will ignore all those complications, and study exclusively the $V_L V_L \to V_L V_L$ subprocess.[⋆] That is, we will ignore the source of the longitudinally polarized vector bosons (although we will comment on this) and, to quote numerical results, we will use the effective W approximation [10] In the same way we will ignore all the detection issues that go into analyzing the decay of the longitudinal vector bosons.

In pp colliders there are three main mechanisms to produce vector boson pairs. The most important one is through light $q\bar{q}$ annihilation. [11] This process gives rise mostly to transversely polarized vector bosons and has thus been typically considered as a background. This means that one thinks that it is the longitudinal component of the W and Z that will be more sensitive to electroweak symmetry breaking, so that one imposes cuts and defines signals in a way that will decrease the relative importance of this mechanism. However, this is not the complete story. There are some longitudinal pairs produced in this way that are always in an $I = 1$ state, when the light quark masses are neglected.[†] This production mechanism (as well as e^+e^- annihilation [12]) is thus very sensitive to "new physics" with a vector resonance like a techni-rho. Moreover, one typically considers the production to be given by the standard model couplings. However, in a more general scenario there are other couplings and vector boson pair production in $q\bar{q}$ annihilation is in principle sensitive to those additional couplings. [13]

⋆ V_L stands for W_L or Z_L. We will also loosely call "isospin" the custodial symmetry relating w and z.

† See the second paper of Ref. 7.

A second mechanism for producing vector boson pairs is gluon fusion. [14] In this case two gluons turn into two vector bosons via an intermediate quark loop. It is only heavy quarks in the loop that give rise to longitudinally polarized vector bosons in a significant amount. In this case, the $V_L V_L$ pair is in an $I = 0$ state,* and thus this channel is particularly sensitive to new physics with a scalar resonance like a heavy Higgs boson. This mechanism could dominate the production of longitudinally polarized vector boson pairs in theories with additional pseudo-Goldstone bosons that carry color and thus couple to gluons. [15]

Finally, there is the vector boson fusion process which is important only in the case of a strongly interacting electroweak symmetry breaking sector. [16] It is thus this mechanism that would provide the cleanest signals if one could isolate it. We will concentrate on this mechanism for our discussion and we will look only at the $V_L V_L \to V_L V_L$ subprocess, having in mind that we can use the effective W approximation if we need numerical estimates.

2. Strongly Interacting Electroweak Symmetry Breaking Sector

In the minimal standard model electroweak symmetry is spontaneously broken by a scalar particle, the Higgs boson. One can study in this case what is the scattering amplitude for longitudinally polarized vector bosons at high energies. If we look at this amplitude in the $I = 0$ channel we find: [17]

$$\mathcal{M} = \frac{1}{v^2}\left(3s + t + u - \frac{3s^2}{s - M_H^2} - \frac{t^2}{t - M_H^2} - \frac{u^2}{u - M_H^2}\right)$$

and if we project out the $J = 0$ partial wave, we obtain for $s \gg M_H^2$:

$$a_0^0 \to \frac{5M_H^2}{32\pi v^2}$$

a result that is proportional to M_H^2. We also know, that a minimal consequence of partial wave unitarity is that

$$|\mathrm{Re}a_0^0| \leq \frac{1}{2}$$

So we can see that for a sufficiently large Higgs boson mass (it turns out to be around 800 GeV), this amplitude will "violate unitarity". Of course the amplitude will not violate

* Actually, it is not exactly an $I = 0$ state because of the mass difference between top and bottom; but this will not affect our discussion.

unitarity, but what this means is that the amplitude is becoming sufficiently large that we cannot trust our perturbative calculation. This has been interpreted to mean that the minimal standard model becomes strongly interacting for a sufficiently heavy Higgs boson. [2] This is the simplest example of a strongly interacting electroweak symmetry breaking sector, and has been used as a model by many people to see what could be expected at the SSC.

Other models that have been studied contain other resonances (with other quantum numbers), or no resonances at all. Typical models with resonances include a rescaling of $\pi\pi$ scattering data, [16] and different unitarization prescriptions applied to different perturbative amplitudes. [18] There have also been studies of an $O(2N)$ model in the large N limit. [19] A model without resonances was introduced by Chanowitz and Gaillard, [16] and consisted of allowing each partial wave (for the first few of them) to grow according to the low energy theorems until it saturated partial wave unitarity, it was then assumed to have that constant maximum value for higher energies.

It is partly this latter model that leads to the "no-lose" corollary. That is to say, if we do not find a light Higgs and electroweak symmetry breaking occurs in some strongly interacting theory of W and Z, the SSC will either find new resonances associated with the new strong interaction, or it will find an enhancement in the $V_L V_L$ scattering amplitudes. [4]

For the remainder of this paper we will concentrate on the latter possibility. That is, we will assume that no new resonances are found.

3. Global Symmetries of the Standard Model

The standard model is based on the gauge group $SU(2)_L \times U(1)_Y$ that is spontaneously broken to $U(1)_Q$. The minimal global symmetry consistent with the gauge symmetry is thus a global $SU(2) \times U(1)$ that is broken spontaneously to $U(1)$. The global symmetry groups can also be larger than this one. For example it could be $SU(2) \times SU(2)$ that breaks down to $SU(2)$. In this case there is a "custodial" $SU(2)$ symmetry and $\rho = 1$ (up to electroweak radiative corrections). If the global symmetry group is larger than this, as is the case in some techni-color models, one is left with pseudo-Goldstone bosons after symmetry breaking. Since presumably these are the lightest states associated with such theories they are then expected to dominate their low energy behavior. Since experimentally $\rho \approx 1$ we will for simplicity consider only the case with a custodial $SU(2)$. The minimal standard model with an elementary Higgs boson has these global symmetries (an $O(4)$ global symmetry is broken to $O(3)$).

At low energies compared with the scale of electroweak symmetry breaking, one has to describe the interactions of the Goldstone bosons (or would be Goldstone bosons) associated with the global symmetry breaking. This is done in a very compact way using chiral Lagrangians. The use of effective chiral Lagrangians also allows one to parameterize these interactions is a model independent way, in terms of a few couplings, and to organize them as an expansion in powers of external momenta. [6] It is thus an appropriate tool to describe the physics of Goldstone bosons at energies below the scale of electroweak symmetry breaking. For $SU(2) \times SU(2) \to SU(2)$, one has exactly the same situation as that of QCD with two massless quarks (u, d), so that all the studies of pion scattering can be applied to longitudinal vector boson scattering by invoking the equivalence theorem.

We can introduce the would-be Goldstone boson fields, w^+, w^-, and z through the matrix $\Sigma = \exp(i\vec{\tau} \cdot \vec{w}/v)$. The lowest order effective Lagrangian describing the scalar sector of a general model for electroweak symmetry breaking contains two derivatives and is unique:

$$\mathcal{L} = \frac{v^2}{4} \mathrm{Tr} \partial_\mu \Sigma \partial^\mu \Sigma^\dagger \tag{3.1}$$

The gauge interactions of the standard model are introduced by requiring the Lagrangian in Eq. (3.1)to be gauge invariant under $SU(3)_c \times SU(2)_L \times U(1)_Y$. This is accomplished by replacing the derivative with a covariant derivative:

$$\partial_\mu \Sigma \to \mathcal{D}_\mu \Sigma = \partial_\mu \Sigma - i\frac{g_S}{2} G^\alpha_\mu [\lambda^\alpha, \Sigma] - i\frac{g}{2} W^\alpha_\mu \tau^\alpha \Sigma + i\frac{g'}{2} B_\mu \Sigma \tau_3 \tag{3.2}$$

If one expands Eq. (3.1) after introducing Eq. (3.2), and looks at the terms quadratic in the gauge fields, one can read that $M_W = gv/2$, so that our choice of normalization in Eq. (3.1) corresponds to the usual $v \approx 246\ GeV$.

We will now ignore the gauge interactions and concentrate in the purely scalar sector of the theory. The next to leading effective Lagrangian contains two free parameters that characterize the underlying theory:

$$\mathcal{L}^{(4)} = \alpha_1 \mathrm{Tr}(\partial_\mu \Sigma \partial^\mu \Sigma^\dagger) \mathrm{Tr}(\partial_\nu \Sigma \partial^\nu \Sigma^\dagger) + \alpha_2 \mathrm{Tr}(\partial_\mu \Sigma \partial_\nu \Sigma^\dagger) \mathrm{Tr}(\partial^\mu \Sigma \partial^\nu \Sigma^\dagger) \tag{3.3}$$

It is worth emphasizing at this point that this is the most general form for the effective Lagrangian at this order only in the absence of gauge interactions.

In order to calculate amplitudes to next to leading order in powers of the external momenta, p^4, one then uses Eq. (3.1) at tree and one-loop levels, and Eq. (3.3) at tree level. The infinities that appear when using Eq. (3.1) at one loop, must have the form of Eq. (3.3), since this is the most general form at order p^4, and can thus be absorbed by defining renormalized parameters α_1 and α_2.

The important point is simply that any theory with a global $SU(2) \times SU(2)$ spontaneously broken to $SU(2)$ will have the same lowest order effective Lagrangian. This means, for example, that the $\mathcal{O}(p^2)$ terms of ww (and by the equivalence theorem of $V_L V_L$) scattering amplitudes are simply given by the low energy theorems for $\pi\pi$ scattering derived by Weinberg [20] many years ago with the replacement $f_\pi \to v$.

At next to leading order one finds corrections to the lowest order behavior of amplitudes, that do depend on the underlying dynamics (which in the effective Lagrangian language is thus parameterized by α_1 and α_2). This description of scattering amplitudes has been seen to work reasonably well for the case of $\pi\pi$ scattering up to energies of about 500 MeV. [21] In that case, the scale of chiral symmetry breaking is about 1 GeV, so that we may naively expect the effective Lagrangian description of ww scattering to be "reasonable" below about 1.5 TeV.* Since this is the range of ww pair invariant masses that the SSC is likely to be able to probe, we can use the formalism to explore (at least qualitatively) aspects of ww scattering in SSC physics.

4. $V_L V_L$ Scattering Amplitudes

The custodial $SU(2)$ symmetry that we are assuming, together with crossing symmetry, relates all the $V_L V_L \to V_L V_L$ scattering amplitudes in such a way that there is only one independent function of the Mandelstam variables, $A(s,t,u)$, needed to construct all the amplitudes. This function is equal to the amplitude for $w^+ w^- \to zz$, and we will thus concentrate on this process.

In the minimal standard model this amplitude is easily computed at tree level to be:

$$
\begin{aligned}
iA(s,t,u) = i\mathcal{M}(w^+ w^- \to zz) &= -i\frac{M_H^2}{v^2} + \left(\frac{-iM_H^2}{v}\right)^2 \frac{i}{s - M_H^2} \\
&= -i\frac{M_H^2}{v^2}\frac{s}{s - M_H^2} \\
&= i\left(\frac{s}{v^2} + \frac{s^2}{v^2 M_H^2} + \cdots\right)
\end{aligned}
\tag{4.1}
$$

In the last line the amplitude has been expanded up to next to leading order in powers

⋆ This is because the scale of electroweak symmetry breaking is expected to be about $4\pi v \sim 3\ TeV$. [22]

of s assuming $s \ll M_H^2$ to illustrate the idea behind the use of the effective Lagrangians. Whereas the first term is just the universal low energy theorem, the second term depends on the Higgs mass, which in this case is the information associated with the symmetry breaking sector. The idea is thus to measure deviations in the scattering amplitudes from their universal lowest order behavior, and to try to relate these to models for the underlying dynamics.

In terms of the effective Lagrangian, Eqs. (3.1), (3.3), the amplitude is given by:

$$\begin{aligned} A(s,t,u) =& \frac{s}{v^2} + \frac{4}{v^4}\Big(2\alpha_1(\mu)s^2 + \alpha_2(\mu)(t^2+u^2)\Big) \\ &+ \frac{1}{16\pi^2 v^4}\Big[-\frac{1}{12}(3t^2+u^2-s^2)\ln\Big(-\frac{t}{\mu^2}\Big) \\ &- \frac{1}{12}(3u^2+t^2-s^2)\ln\Big(-\frac{u}{\mu^2}\Big) - \frac{s^2}{2}\ln\Big(-\frac{s}{\mu^2}\Big)\Big] \end{aligned}$$

The first term comes from the tree level lowest order Lagrangian, the second term from the tree level $\mathcal{O}(p^4)$, Eq. (3.3) Lagrangian, and the last term is the one loop amplitude with the vertices from Eq. (3.1). The scale dependence of this last term is compensated by the scale dependence of α_i in physical amplitudes.

If one has a specific model for electroweak symmetry breaking, one can calculate the amplitudes directly. Such a calculation has been done to one loop in the case of the minimal standard model. [23] † If one expands that result for low energies one can read off the values of the couplings α_i in that model:

$$\begin{aligned} \alpha_1(\mu) =& \frac{v^2}{8M_H^2} + \frac{1}{4}\Big[-\frac{1}{16\pi^2}\Big(\frac{9\pi}{4\sqrt{3}} - \frac{37}{9}\Big) - \frac{1}{48\pi^2}\ln\Big(\frac{\mu}{M_H}\Big)\Big] \\ \alpha_2(\mu) =& \frac{1}{4}\Big[-\frac{1}{16\pi^2}\Big(\frac{2}{9}\Big) - \frac{2}{48\pi^2}\ln\Big(\frac{\mu}{M_H}\Big)\Big] \end{aligned}$$

The first term in α_1 comes from tree level Higgs exchange, whereas the other terms arise from one loop diagrams. The coefficients of the logarithms of those terms are fixed, but the constant factors depend on the renormalization prescription. As given above, they correspond to the renormalization of Ref. 23.

† Since there appears to be some confusion in the literature we must emphasize that these calculations were done using ordinary perturbation theory in the coupling constant λ, which is proportional to M_H^2, and are thus valid **only** for a "light" Higgs boson.

Clearly we cannot calculate the couplings α_i for an unknown underlying theory. However, it is possible to relate these couplings to the properties of nearby resonances.[24] To do so one couples the Goldstone boson fields to resonances of the desired quantum numbers.[25] For example, the general coupling to a scalar, isoscalar resonance H is of the form:

$$\mathcal{L}^S = g_H H \mathrm{Tr} \partial_\mu \Sigma \partial^\mu \Sigma^\dagger \tag{4.2}$$

and the coupling to a vector, isovector resonance has the form

$$\mathcal{L}^\rho = g_\rho \rho^i_{\mu\nu} \mathrm{Tr}\left(\tau^i \xi^\dagger \partial^\mu \Sigma \partial^\nu \Sigma^\dagger \xi \right) \tag{4.3}$$

where we have used an antisymmetric tensor field for the resonance, and we have introduced $\xi = \exp i\vec{\tau} \cdot \vec{w}/2v$.

With these couplings one can compute the resonant ww scattering amplitudes and expand them to order p^4 to obtain equivalent values for α_i.* Doing this, and rewriting the couplings in terms of the resonance width one finds:‡

$$\begin{aligned} \alpha_1^H &= \frac{4\pi}{3} \frac{\Gamma_H v^4}{M_H^5} \\ \alpha_2^H &= 0 \end{aligned}$$

in agreement with Eq. (4.1), since in the perturbative regime the Higgs width in the standard model is $\Gamma_H = \frac{3M_H^3}{32\pi v^2}$. One also finds:

$$\begin{aligned} \alpha_1^\rho &= -12\pi \frac{\Gamma_\rho v^4}{M_\rho^5} \\ \alpha_2^\rho &= -\alpha_1^\rho \end{aligned}$$

These simple results correspond to the low energy effect of a tree level exchange of the resonance: "resonant scattering". They will be modified by Goldstone boson loops among other things. In a schematic notation we can write:

$$\alpha_i(\mu) = \sum_R \alpha_i^R + \alpha_i^B(\mu)$$

where we separate the tree level resonant scattering from the "background", such as heavy fermion loops[26] , Goldstone boson loops (which will generate the μ dependence) and

* More generally one speaks of "integrating out" the resonance field to produce an effective action that is then expanded in powers of external momenta.

‡ See for example the first paper of Ref. 7.

others. Our question can then be rephrased as follows: Is it possible to infer the presence (and perhaps mass and width) of a resonance by studying Goldstone boson scattering at energies below threshold for production of this resonance? As one can see from the above formulas the effect of a resonance on low energy scattering becomes more important for wider and lighter resonances as one would naively expect. For very heavy or narrow resonances the low energy scattering will be dominated by the "background".

A simple exercise consists of taking various values for the mass and width of resonances and seeing if "low energy" data at the SSC could distinguish between them. This was done in the second paper of Ref. 7, with somewhat pessimistic conclusions.♮ In Fig. 1 we show the predictions of the low energy theorems for two channels (solid line). We also show these lowest order amplitudes plus the leading corrections from α_i for a vector resonance with several values for its mass and width; we include one with mass 2 TeV and width $\sim 0.4\ TeV$ which is a scaled up version of QCD (dashed lines). The case of a "standard model" Higgs of mass 2 TeV is also shown for comparison (dotted line), although perturbation theory clearly breaks down for such a heavy Higgs (if we took things literally, this "Higgs boson" would have a width of $\sim 4\ TeV$). The band between dotted lines includes values between $2-3\ TeV$ for the scalar mass and several choices for widths ranging from $M_H/2$ to $2\ M_H$. One can see how at 1 TeV the corrections to the low energy theorems are sizable, and the two models give different amplitudes. The parameters corresponding to a scalar resonance tend to enhance the low energy theorem predictions for these two channels, while those corresponding to a vector resonance do not. Unfortunately, these two channels have large backgrounds (not shown in the figure). In the two channels that are not shown, the two regions actually overlap, indicating that a measurement of deviations from the low energy theorems in those channels would not place very strong constrains on any "new physics".

Finally we turn to the second question we wanted to answer by looking for a "worse case scenario" in which there are no resonances (or they are pushed to high energy scales where they do not affect the scattering in the few TeV energy region). In Fig. 2 we illustrate the model without resonances of Chanowitz and Gaillard[16] . In this case the partial wave amplitudes are allowed to grow with the low energy theorems until they reach a unitarity bound, after which they are allowed to saturate. We also show our model[8] where we vary α_1 and α_2 and choose them in such a way as to postpone the violation of

♮ We refer the reader to that reference for a complete set of curves for all channels.

the unitarity condition as much as possible. We find that with:

$$\begin{aligned} \alpha_1 =& -0.00167 \\ \alpha_2 =& 0.00147 \end{aligned} \tag{4.4}$$

violation of the condition $|\text{Re}T^I_J| \leq 1/2$ is delayed to 2 TeV where the T^1_1 amplitude violates it. Since this energy scale is already outside of the reach of the SSC we use these parameters as a "worse case scenario" where the resonances have been pushed up in energy as much as possible.

In Fig. 3 we compare the results of our model with those of Chanowitz and Gaillard. The results of the two models are remarkably similar for W^+W^- and W^+Z final states (at least below 1.5 TeV). However, the model with parameters from Eq. (4.4) gives somewhat lower amplitudes in the ZZ and W^+W^+ channels. This is important since it has been argued in the literature that the latter is the channel with the lowest backgrounds.

With our model, we estimate that a signal involving only electrons or muons, in the W^+W^+ channel (where the background is about 3.5 events per SSC year), would consist of about 5 events per SSC year (assuming an annual luminosity of 10 fb^{-1}). For a more precise definition of the signal, and references to background calculations we refer the reader to Ref. 8. Here, we just want to illustrate the point that one is talking about signals that would be very difficult to detect. However, there is some disagreement in the literature about the precise way to arrive at these numbers.[27] This involves both theoretical issues (as details of the structure functions one should use) and more experimental issues (like what are the best signals and the best cuts to get rid of backgrounds). The resolution of these important questions will ultimately determine whether there is really a "no-lose corollary" or not.

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FIGURE CAPTIONS

1. We show for two channels the predictions of the low energy theorems: solid line. We also show in the region between dashed lines the $\mathcal{O}(s^2)$ amplitudes in vector dominated theories with values of M_ρ and Γ_ρ that include a scaled up version of QCD. In the region between dotted lines we show the $\mathcal{O}(s^2)$ amplitudes for scalar dominated theories including a 2 TeV Higgs with a width as predicted in the standard model.

2. We show schematically for a partial wave the model without resonances of Chanowitz and Gaillard and our "worse case scenario".

3. In this figure we compare the subprocess cross sections for the two models without resonances. The dashed line corresponds to the model of Chanowitz and Gaillard, the solid line to our model. The curves (a), (b), (c) and (d) correspond to the final states W^+W^-, ZZ, W^+Z, and W^+W^+.

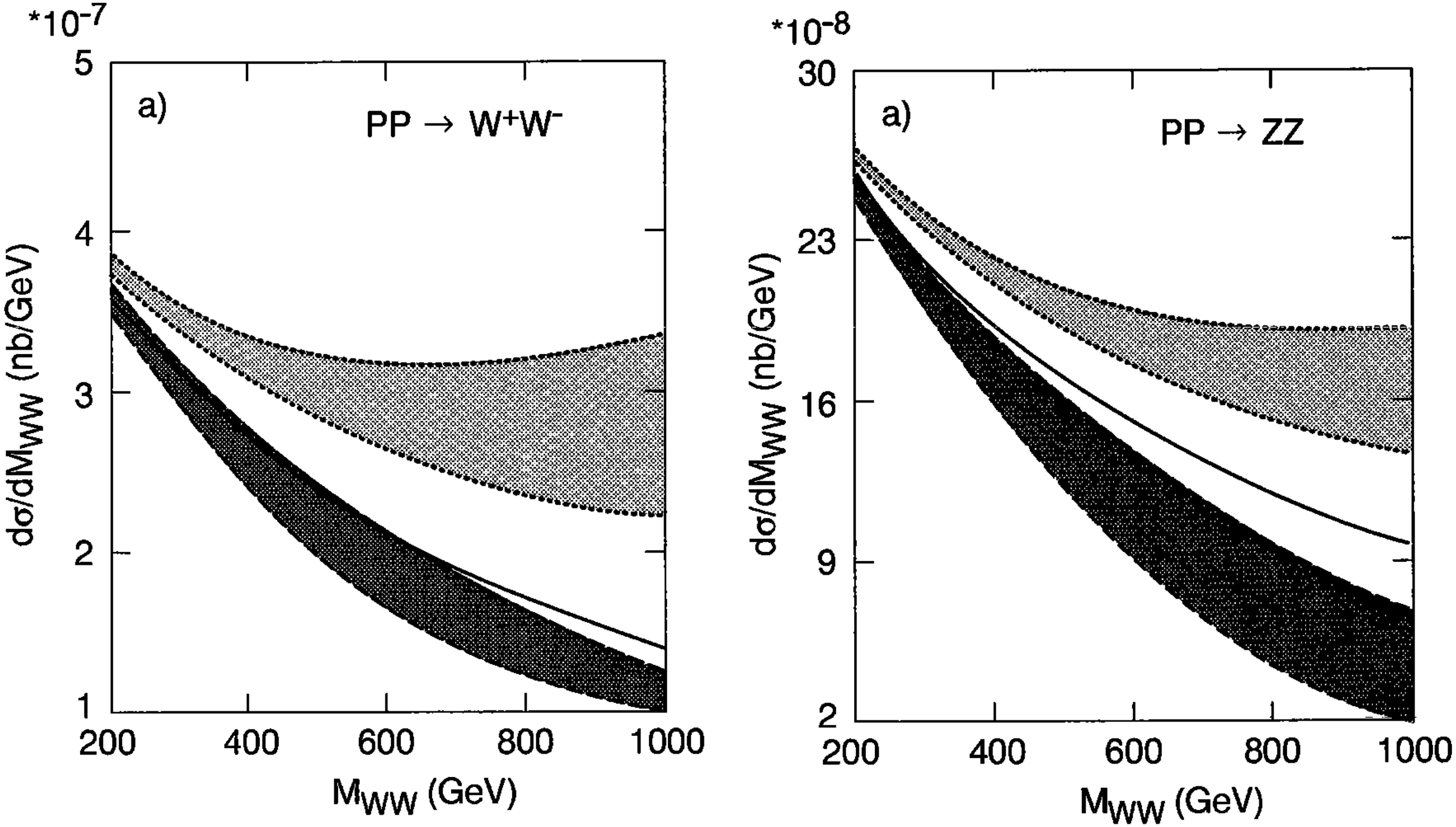
*10-7
5
4
3
2
1
a)
PP → W+W-
dσ/dMWW (nb/GeV)
200
400
600
800
1000
MWW (GeV)
*10-8
30
23
16
9
2
a)
PP → ZZ
dσ/dMWW (nb/GeV)
200
400
600
800
1000
MWW (GeV)

Fig. 1

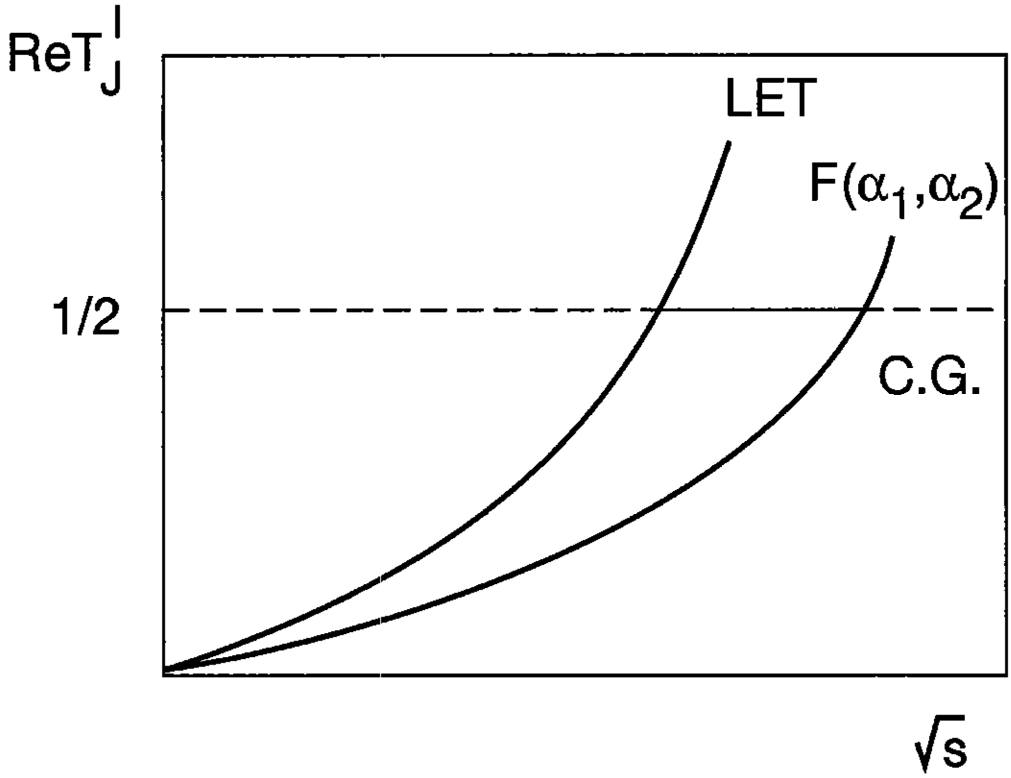

Fig. 2

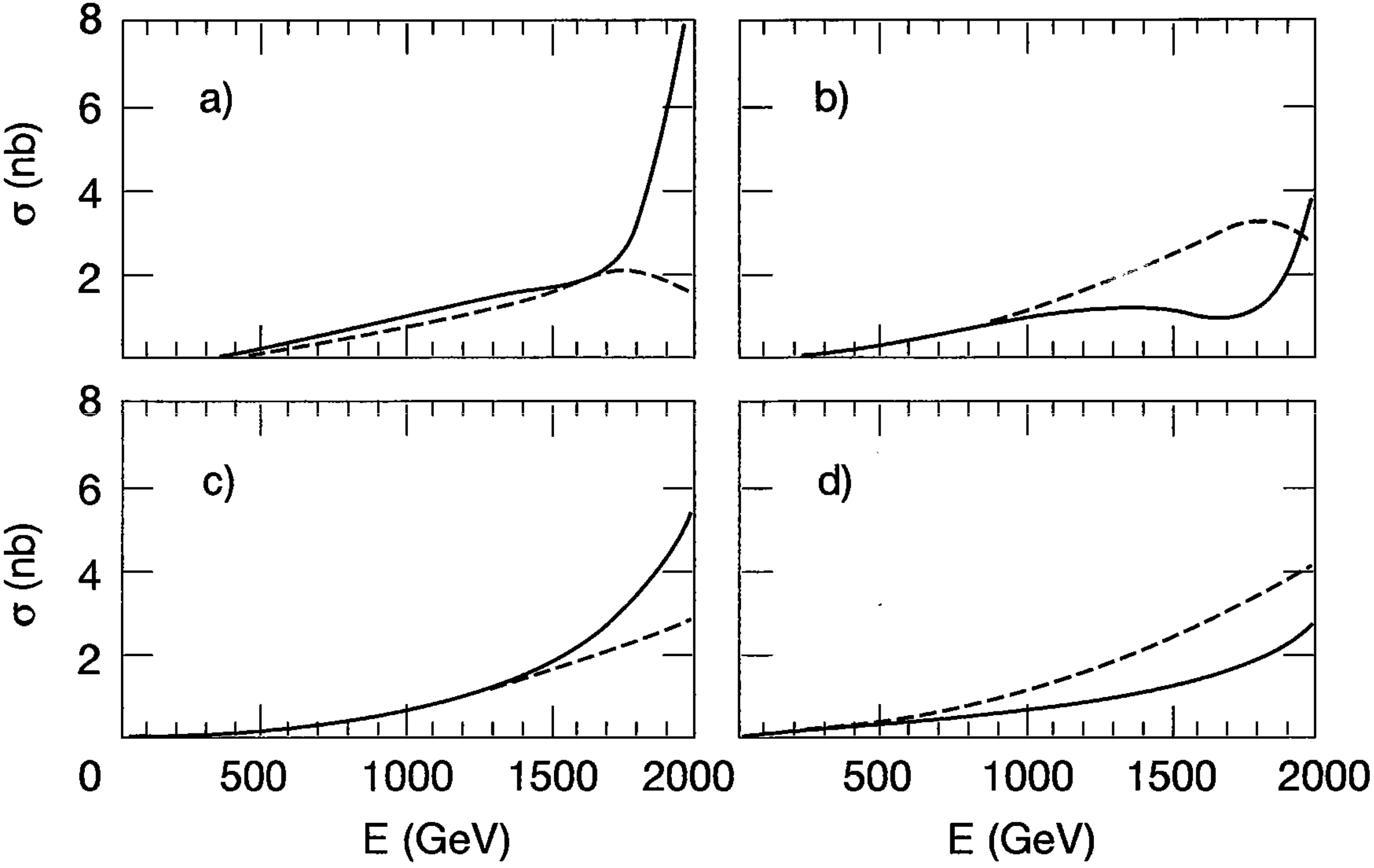
a)
b)
c)
d)
σ (nb)
σ (nb)
8
6
4
2
8
6
4
2
0
500
1000
1500
2000
500
1000
1500
2000
E (GeV)
E (GeV)

Fig. 3