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TESTING ELECTROWEAK SYMMETRY BREAKING THROUGH GLUON FUSION AT pp COLLIDERS

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Abstract

We use chiral Lagrangians to study the nonresonant production of longitudinally polarized vector bosons through gluon fusion at pp colliders. We compute the contributions induced by loops of colored pseudo-Goldstone bosons and colored fermions. We find that the resulting cross-sections potentially dominate the standard-model predictions and provide an important probe of the electroweak symmetry-breaking sector.

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In the standard model of Glashow, Salam and Weinberg, the electroweak symmetry is broken by a set of fundamental scalars called a Higgs doublet. During the past decade, many features of this model have been carefully tested, but no experimental evidence for the Higgs has emerged. While this is not yet a cause for alarm, it leads one to speculate that the Higgs sector might be a rich source of new physics.

On general grounds, we know that whatever replaces the Higgs must give mass to the $W^{\pm}$ and the Z. This implies that the new physics must give rise to three Goldstone bosons, $w^{\pm}$ and z, which become the longitudinal components of the massive vector bosons. Therefore we expect that the interactions of the longitudinal vector bosons will shed light on the dynamical mechanisms which underlie the electroweak symmetry breaking.

Chiral Lagrangians provide a particularly useful framework for studying longitudinal vector bosons because they allow a systematic and model-independent approach to electroweak symmetry breaking. They have previously been used to compute signals and backgrounds for new resonances that might appear at the SSC and LHC. It is equally important, however, to find nonresonant signals of new physics, for planned colliders might not have sufficient reach to reveal the resonant structures.[1]

In this letter we will use chiral Lagrangians to compute longitudinal vector production via gluon fusion. This process has the potential to dominate the production of $Z_L Z_L$ and $W_L^+ W_L^-$ pairs when the electroweak symmetry breaking involves colored particles. Such particles typically arise when one considers the large global symmetry groups associated with realistic theories of dynamical symmetry breaking. In particular, these theories tend to have a large number of relatively light colored pseudo-Goldstone bosons which dramatically enhance the $gg \rightarrow Z_L Z_L$ and $gg \rightarrow W_L^+ W_L^-$ production amplitudes.[2] Gluon fusion provides an important tool for testing the physics that breaks the electroweak symmetry.

Before computing the scattering amplitudes, we must first construct the chiral Lagrangian that describes the relevant couplings. We make the simplifying assumption that whatever physics breaks the electroweak symmetry, it has a global chiral symmetry group $G_L \times G_R$, spontaneously broken to the diagonal subgroup G. The *dim* G Goldstone bosons Π^A can then be parametrized by a matrix $\Sigma = \exp(2i\Pi^A T^A/v)$, where the T^A are the generators of G (normalized so $\mathrm{Tr}\, T^A T^B = \frac{1}{2}\delta^{AB}$), and $v = 246$ GeV characterizes the scale of the symmetry breaking. Of course, three of the Goldstone bosons are exactly massless and become the longitudinal components of the $W^\pm$ and the Z. The remaining *dim* $G - 3$ Goldstones must acquire mass; they are known as pseudo-Goldstone bosons.

Under a global chiral transformation, the matrix Σ transforms as $\Sigma \to L\Sigma R^{-1}$, where $L, R \in G$. For this model to describe electroweak symmetry breaking, the $SU(3) \times SU(2) \times U(1)$ gauge group must be embedded in the chiral symmetry group $G_L \times G_R$. To define the embeddings, we first construct the matrices $X^\alpha = X^{\alpha A} T^A$, $X^a = X^{aA} T^A$ and $X = X^A T^A$, which generate $SU(3)$, $SU(2)$ and $U(1)$, respectively. These matrices are in the Lie algebra of G and are normalized in the same way as the generators T^A. (Note, however, that $\mathrm{Tr}\, X X^a = \frac{1}{2}\,\delta^{a3}$.) The embedding is then defined by the covariant derivative

$$D_\mu \Sigma \;=\; \partial_\mu \Sigma \;-\; i g_s G^\alpha_\mu [X^\alpha, \Sigma] \;-\; i g W^a_\mu X^a \Sigma \;+\; i g' B_\mu \Sigma X \;, \qquad (01)$$

where G^α_μ is the $SU(3)$ color gauge field, and W^a_μ and B_μ are the the $SU(2) \times U(1)$ gauge bosons. Note that the gauge couplings explicitly break the global symmetry group $G_L \times G_R$.

The self-interactions of Goldstone bosons and their interactions with the standard model gauge fields are given by the chiral Lagrangian. The Lagrangian is nonrenormalizable, but it makes sense as an effective theory for energies $s \lesssim \Lambda^2$, where $\Lambda \sim 1$ TeV denotes the scale of the physics responsible for breaking the electroweak symmetry. To lowest order in the energy expansion, the effective

Lagrangian is given by

$$\mathcal{L} \;=\; \frac{v^2}{4}\,\mathrm{Tr}\,D_\mu\Sigma D^\mu\Sigma^\dagger \;-\; \mathrm{Tr}\;\Sigma M\Sigma^\dagger M^\dagger \;. \tag{02}$$

The first term is completely determined by v and the embedding matrices X. Expanding in powers of the Goldstone fields Π^A, one can see that it gives mass to the $W^\pm$ and the Z, with $M_W = M_Z\cos\theta$ in accord with experimental measurements. To higher order, the term contains other couplings of interest, including the self-couplings of the Goldstone bosons,

$$\mathcal{L}^{\Pi\Pi\Pi\Pi} \;=\; -\frac{1}{6v^2}\,f^{ABE}f^{CDE}\,\Pi^A\partial_\mu\Pi^B\Pi^C\partial^\mu\Pi^D \;, \tag{03}$$

as well as the couplings of one and two gluons to two Goldstones,

$$\begin{aligned} \mathcal{L}^{G\Pi\Pi} &= g_s\,f^{ABC}\,G^\alpha_\mu\,X^{\alpha A}\,\Pi^B\partial^\mu\Pi^C \\ \mathcal{L}^{GG\Pi\Pi} &= \frac{1}{2}\,g_s^2\,f^{ABE}f^{CDE}\,G^\alpha_\mu G^{\gamma\mu}\,X^{\alpha A}X^{\gamma C}\,\Pi^B\Pi^D \;, \end{aligned} \tag{04}$$

where the f^{ABC} are the structure constants of G.

The second term in (02) represents a contribution to the mass of the pseudo-Goldstone bosons. Gauge invariance requires that M must commute with the matrices X. In general, however, it breaks the rest of the chiral symmetries. The details are very model-dependent, so we will restrict our attention to the massless case, and only plot cross sections in the region $s \gtrsim 4\mu^2$, where μ is a typical eigenvalue of M.

With the Lagrangian (2), we have what we need to compute the amplitudes induced by a loop of colored pseudos. For comparison, we will also compute the contribution from a loop of colored fermions. As above, the couplings are given in terms of a chiral Lagrangian. For simplicity, we assume the fermion couplings preserve parity. This implies that the left-handed fermions ψ_L and the right-handed fermions ψ_R transform under the same representation D of G, $\psi_L \to D(L)\,\psi_L$, $\psi_R \to D(R)\,\psi_R$. The leading-order effective Lagrangian is just

$$\mathcal{L} \;=\; i\bar\psi\gamma^\mu\partial_\mu\psi - \kappa\bar\psi_L\gamma^\mu J_{L\mu}\psi_L - \kappa\bar\psi_R\gamma^\mu J_{R\mu}\psi_R - m\bar\psi_L D(\Sigma)\psi_R - m\bar\psi_R D(\Sigma^\dagger)\psi_L\ , \tag{05}$$

where $\kappa = 1 - g_A$, $J_{L\mu} = i\Sigma D_\mu\Sigma^\dagger$, $J_{R\mu} = i\Sigma^\dagger D_\mu\Sigma$, and the fermion mass m can be different for each irreducible representation in D. Expanding in the Goldstone fields, we find

$$\begin{aligned} \mathcal{L}^{\psi\psi\Pi} &= \frac{2\kappa}{v}\,\bar\psi\gamma^\mu\gamma^5 t^A\psi\,\partial_\mu\Pi^A \;-\; \frac{2im}{v}\,\bar\psi\gamma^5 t^A\psi\,\Pi^A \\ \mathcal{L}^{\psi\psi\Pi\Pi} &= \frac{4\kappa}{v^2}\,f^{ABC}\,\bar\psi\gamma^\mu t^A\psi\,\Pi^B\partial_\mu\Pi^C \;+\; \frac{2m}{v^2}\,\bar\psi t^A t^B\psi\,\Pi^A\Pi^B\ , \end{aligned} \tag{06}$$

where the t^A are the generators of G in the representation D, such that $\mathrm{Tr}\,t^At^B = \frac{1}{2}\delta^{AB}$. The fermion couplings are determined in terms of κ and the mass m. We now have what we need to compute the amplitudes for $gg \to Z_LZ_L$ and $gg \to W_L^+W_L^-$.

Since the global symmetry group is always larger than $SU(2) \times SU(2)$, the effective Lagrangian has a custodial $SU(2)$ symmetry which guarantees that the two amplitudes are identical, $\mathcal{M}(gg \to Z_LZ_L) = \mathcal{M}(gg \to W_L^+W_L^-)$. We shall further simplify our computations by invoking the electroweak equivalence theorem, which states that at energies $s \gg M_W^2$, the scattering amplitudes of longitudinal vectors are equivalent to scattering amplitudes where the longitudinal vectors are replaced by the corresponding would-be Goldstone bosons.[3]

Although the formalism we have presented is completely general, for definiteness we shall present our results in terms of the Farhi-Susskind model of technicolor.[4] This model provides a concrete realization of electroweak dynamical symmetry breaking. In this model, the group $G = SU(8)$, so there are 63 Goldstone bosons. There are also 8 Dirac techni-fermions. The embeddings are such that there is one color octet and two color triplets of weak-triplet pseudo-Goldstone bosons, as well as one color triplet of weak-doublet Dirac techni-fermions. We expect the general features of our numerical results to persist in other models of electroweak dynamical symmetry breaking.

Therefore, in what follows, we shall compute the amplitude for $g^\alpha_\mu(q_1)g^\beta_\nu(q_2) \to z(p_1)z(p_2)$, with the embeddings taken to be those of the Farhi-Susskind model. Gauge invariance implies that the amplitude $\mathcal{M} = \epsilon^\mu(q_1)\epsilon^\nu(q_2)\delta_{\alpha\beta}\mathcal{M}_{\mu\nu}$ must be of the form

$$\begin{aligned} \mathcal{M}_{\mu\nu} &= A(s,t,u)\left(-\frac{s}{2}g_{\mu\nu} + q_{2\mu}q_{1\nu}\right) \\ &+ B(s,t,u)\left(-\frac{ut}{2}g_{\mu\nu} - sp_{1\mu}p_{1\nu} - tq_{2\mu}p_{1\nu} - up_{1\mu}q_{1\nu}\right), \end{aligned} \tag{07}$$

where s, t, and u are the subprocess Mandelstam variables.

We shall first consider the amplitude induced by the (massless) colored pseudo-Goldstone bosons. It is not hard to see that the amplitude vanishes at tree level. The one-loop amplitude is finite, and is given by

$$\begin{aligned} A(s,t,u) &= \sum_R T(R)\left(\frac{\alpha_s}{\pi v^2}\right) \\ B(s,t,u) &= 0\ , \end{aligned} \tag{08}$$

where $T(R) = 1/2$ for each color triplet and $T(R) = 3$ for each color octet.

The contribution from the colored fermions can be readily computed as well.

For simplicity, we take $\kappa = 0$ (or $g_A = 1$). Then the one-loop result is given by[5]

$$\begin{aligned}
A(s,t,u) \;&=\; \left(\frac{\alpha_s}{4\pi m^2 v^2}\right)\left\{\int\limits_0^1 y dy\left(\left[4s\frac{m^2}{t(1-y)}\left(1-\frac{m^2}{ty(1-y)}\right)\right.\right.\right. \\
&\quad + \left. 2u\left(1-\frac{m^2}{ty(1-y)}\right) \;+\; \frac{2m^2}{y(1-y)}\right]f(s,u,t) \;-\; \left.\frac{4}{1-y}\frac{m^4}{u}\right) \\
&\quad + \int\limits_0^1 dy\left[\left(4t+\frac{2u(1-2y)}{1-y}\right)\left(\frac{m^2}{t}f(u,s,t) \;-\; \frac{m^2}{s}f(u,t,s)\right)\right. \\
&\quad + \left.\left. 2tyf(u,t,s) \;-\; 2sy(2y-1)f(u,s,t)\right]\right\} \;+\; (t\leftrightarrow u)
\end{aligned} \tag{09}$$

and

$$\begin{aligned}
B(s,t,u) \;&=\; \left(\frac{\alpha_s}{\pi v^2 m^2}\right)\left\{\int\limits_0^1 dy\; y(y-1)\left(\Big[f(s,u,t)+f(u,t,s)\Big]\right.\right. \\
&\quad + \frac{m^2}{t(1-y)}\left[\frac{m^2}{t(1-y)}-1\right]f(u,s,t) \\
&\quad + \frac{m^2}{s}\left[1-\frac{y}{1-y}+\frac{m^2}{s(1-y)^2}\right]f(u,t,s) \\
&\quad - \left.\left.\frac{y}{1-y}\frac{m^4}{st}\right)\right\} \;+\; (t\leftrightarrow u)\;,
\end{aligned} \tag{010}$$

where

$$f(s,t,u) \;=\; \frac{m^4}{sm^2+uty(1-y)}\log\left(1-\frac{u}{m^2}y(1-y)\right)\;. \tag{011}$$

Note that the contributions to $\mathcal{M}$ from (09) and (010) vanish as s/m^2 for $s \ll m$. They contribute to a local operator in the chiral Lagrangian of order s^2/Λ^4. This is a natural consequence of the fact that the fermion mass term preserves the chiral symmetry. The fermions can be decoupled, and all their effects can be written in terms of local operators in the effective Lagrangian. In contrast, the contribution from the pseudo-Goldstone loop is of order s/Λ^2, and it cannot be written as a gauge-invariant local operator. The pseudo-Goldstone bosons cannot be decoupled, and their effects cannot be absorbed into local operators in the chiral Lagrangian.[6]

We now have what we need to compute the invariant mass distribution for longitudinally polarized Z_L pair production at the SSC and LHC.[7] In Fig. 1 we show the contribution of a color octet, weak triplet of pseudo-Goldstone bosons to the process[8] $pp \to gg \to Z_L Z_L$. The contribution of a color triplet is 1/36 of that of the octet. For comparison, we also show the contribution to Z pair production from quark-antiquark annihilation. The $q\bar{q}$ curve includes both the transverse and longitudinal polarizations of the Z's.

In Fig. 1 we see that the contribution of colored pseudo-Goldstones to Z pair production can be very large. In the Farhi-Susskind model, it dominates the standard model contribution for any Higgs mass. This results from the large octet color factor and from the fact that the gluon luminosity is much larger than that for quarks. We did not include a mass for the pseudo-Goldstones, so the figure is valid only for $M_{ZZ}^2 \gtrsim 4\mu^2$, where μ is the pseudo-Goldstone mass. Below threshold, the curves depend on the mass matrix M.

In Fig. 2 we plot the differential cross section for $pp \to gg \to Z_L Z_L$ from a mass-degenerate color triplet of weak-doublet fermions. We take $\kappa = 0$ and $m = 500$ GeV. For comparison, we also show the contribution from a 200 GeV top quark.[9] As expected, the contributions of heavy fermions are significantly smaller than those from colored pseudo-Goldstone bosons. This follows from the fact that the fermion couplings are suppressed in the energy expansion by one extra factor of s.

In this letter we have computed the contributions to $Z_L Z_L$ and $W_L^+ W_L^-$ production induced by the colored pseudo-Goldstone bosons that arise in models of dynamical electroweak symmetry breaking. At future hadron colliders such as the SSC or LHC, these contributions are potentially large, and are typically more important than the contribution from the top quark or from new generations of quarks or techniquarks. The observation of an enhanced gauge boson pair production would provide important insights into the mechanism for electroweak symmetry breaking.

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2. Of course, if such particles exist, they would be directly produced at the SSC or LHC. Their decays, however, are very model-dependent. We are making the more model-independent statement that colored pseudo-Goldstone bosons also contribute significantly to $Z_L Z_L$ and $W_L^+ W_L^-$ production.
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4. E. Farhi and L. Susskind, *Phys. Rev.* **D20** (1979) 3404; *Phys. Rep.* **74** (1981) 277.
5. This expression also holds for a standard-model quark doublet in the limit of infinite Higgs mass.

6. A similar argument implies that contributions from vector resonances occur at the same order as those from fermions, and that they decouple with mass as well.

7. We always use the EHLQ structure functions with $\Lambda = 200$ MeV. See E. Eichten, I. Hinchliffe, K. Lane and C. Quigg, *Rev. Mod. Phys.* **56** (1984) 1.

8. We remind the reader that because of the custodial $SU(2)$ symmetry, the amplitudes for $Z_L Z_L$ and $W_L^+ W_L^-$ production are identical. There is, however, an extra factor of 1/2 in the $Z_L Z_L$ cross section because of Bose symmetry.

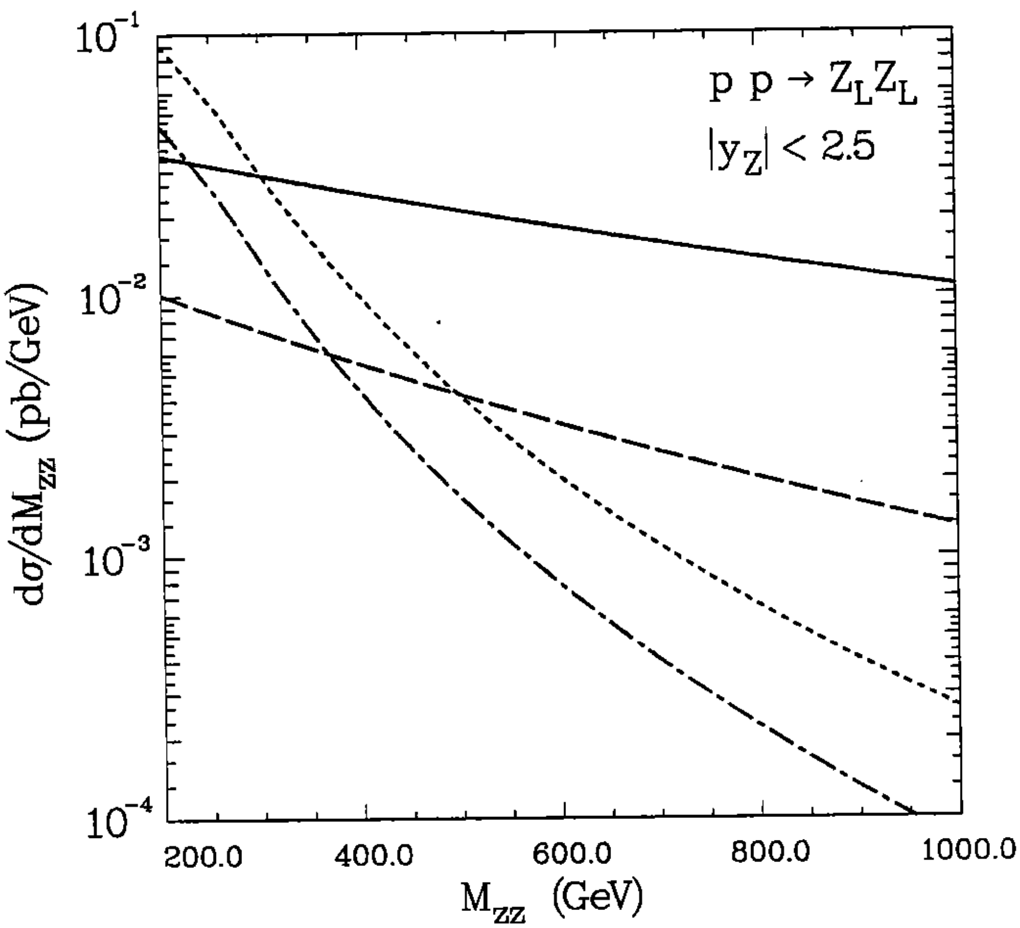
9. The full amplitude for $pp \to gg \to ZZ$ has been computed in the standard model by E. Glover and J. Van der Bij, *Nucl. Phys.* **B321** (1989) 561. Our results for the top quark contribution to $Z_L Z_L$ pair production agree with theirs.

FIGURE CAPTIONS

1. $d\sigma/dM_{ZZ}$ for pp collisions with $|y_Z| < 2.5$. The solid (dashed) curve gives the contribution to $gg \to Z_L Z_L$ from a loop of color octet, weak-triplet pseudo-Goldstone bosons at $\sqrt{S} = 40$ TeV ($\sqrt{S} = 16$ TeV). The dotted (dot-dashed) line shows the contribution from $q\,\overline{q} \to ZZ$ at $\sqrt{S} = 40$ TeV ($\sqrt{S} = 16$ TeV).

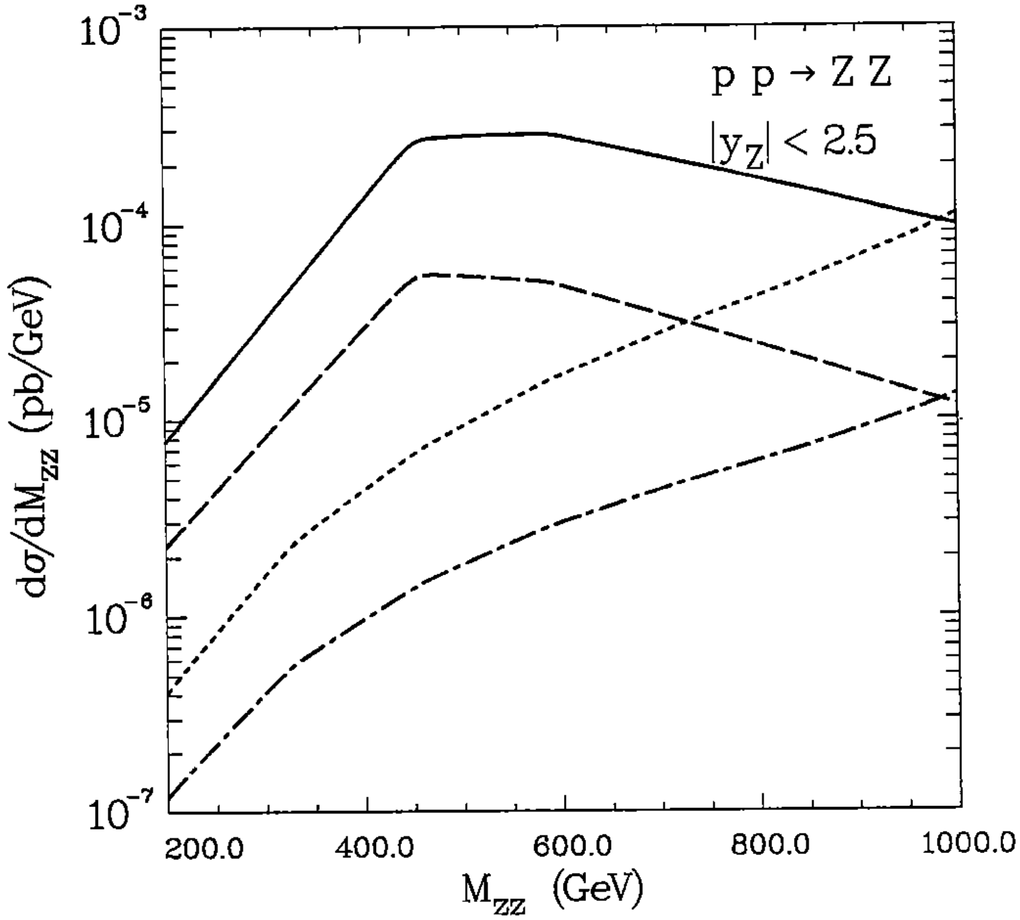
2. $d\sigma/dM_{ZZ}$ in pp collisions with $|y_Z| < 2.5$. The solid (dashed) line gives the contribution of a 200 GeV top quark to $gg \to Z_L Z_L$ at $\sqrt{S} = 40$ TeV ($\sqrt{S} = 16$ TeV). The dotted (dot-dashed) line gives the contribution of a color triplet, weak doublet of heavy fermions with $m = 500$ GeV at $\sqrt{S} = 40$ TeV ($\sqrt{S} = 16$ TeV).

Fig. 1



p p → $Z_L Z_L$
$|y_Z| < 2.5$
10^{-1}
10^{-2}
10^{-3}
10^{-4}
$d\sigma/dM_{ZZ}$ (pb/GeV)
200.0
400.0
600.0
800.0
1000.0
M_{ZZ} (GeV)

Fig. 2



p p → Z Z
|y_Z| < 2.5
10^-3
10^-4
10^-5
10^-6
10^-7
dσ/dM_ZZ (pb/GeV)
200.0
400.0
600.0
800.0
1000.0
M_ZZ (GeV)