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ABSTRACT

We study the production of heavy leptons belonging to a fourth generation,through the vector boson fusion mechanism in e^+e^- collisions at CLIC energies. Lepton production through photon fusion is a more efficient mechanism since the total cross section is larger for a considerable range of lepton masses. Only for very high masses the fusion of longitudinally polarized bosons becomes competitive with photon fusion.

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I. INTRODUCTION

The standard model of the electroweak and strong interactions(1,2) has met remarkable experimental success in the last years. However, there are fundamental questions that are beyond the scope of the standard model whose answers might be within the reach of the future generation of accelerators. One of these questions is the number of fermionic families. The intriguing replication of leptons (and quarks) has led to a plentiful set of models where the existence of a fourth family is admitted(3). We can also recall that some superstring models predict an even number of families, therefore we might find at least one extra lepton(4).

Admitting the existence of a fourth heavy lepton, we have recently shown that the most efficient process for the production of these leptons in hadronic colliders is the vector–boson–fusion mechanism(5). This result was obtained within the effective–vector–boson approximation(6,7). As happens for the Higgs boson sector, the physics of high energy colliders may be dominated by vector–boson–fusion processes(8), and many of these have been studied recently in the case of hadronic colliders(9), but not quite so extensively for e^+e^- machines.

In this paper we study the production of heavy leptons via vector–boson–fusion at CLIC energies. This machine is in study at CERN and is proposed to operate at $\sqrt{s} = 2$ TeV with a luminosity $\mathscr{L} = 10^{34}$ cm^{-2}s^{-1} (or an integrated luminosity $L = 10^5$ pb/year). We shall compare the cross sections of photons and weak (longitudinally and transversally polarized) vector–bosons–fusion giving a pair of charged leptons, or a heavy lepton and its neutrino, which will be computed within the effective–vector–boson approximation. Section II contains a discussion of this quoted approximation and display the basic tools we shall deal with. In section III the analytical cross sections of each subprocess can be found, and we leave for section IV our results and conclusions.

II. THE EFFECTIVE–VECTOR–BOSON APPROXIMATION

The effective–vector–boson approximation is an extension to massive weak gauge bosons of the Weizsäcker–Williams (or leading logarithmic) approximation. It is well known that this method leads to quite good results for the two–photon process, i.e., $e^+e^- \to e^+e^-\gamma\gamma \to e^+e^-X$, where in the limit of parallel momenta the photon distribution inside the electron is

$$A_f = \frac{\alpha}{2\pi} \frac{1 + (1-x)^2}{x} \ln\left[\frac{\hat{s}}{4m_e^2}\right] \quad , \tag{1}$$

and the full cross section for the two–photon process is the product of the photon distributions times the cross section of the subprocess $\gamma\gamma \to X$.

The above procedure was generalized, in the case of weak vector bosons, by Kane et al. and Dawson$^{(6)}$, who determined the vector boson distributions inside a fermion for longitudinally and transversally polarized bosons. In the leading order, they are respectively given by

$$V_f^L \cong \frac{\alpha}{\pi} (C_V^2 + C_A^2) \frac{1-x}{x} \quad , \tag{2a}$$

$$V_f^T \cong \frac{\alpha}{2\pi} (C_V^2 + C_A^2) \frac{1 + (1-x)^2}{x} \ln\left[\frac{\hat{s}}{M_V^2}\right] \quad , \tag{2b}$$

where M_V is the gauge boson mass, $\sqrt{\hat{s}}$ the subprocess invariant mass, and

$$C_V = -C_A = \frac{1}{2\sqrt{2}\ \sin\theta_W}$$

for charged ($W^{\pm}$) weak bosons and

$$C_V = \frac{1}{\sin\theta_W\ \cos\theta_W} \left[\frac{1}{2} T_3 - Q \sin^2\theta_W\right]$$

$$C_A = -\frac{1}{2\sin\theta_W\ \cos\theta_W} T_3$$

for the Z^0 boson.

The cross section for $e^+e^- \to V_iV_j \to X$, where $V_{i(j)}$ is any of the electroweak bosons, can be written as

$$\sigma(e^+e^- \to V_iV_j \to X) = \frac{1}{1+\delta_{ij}} \int_{\tau_m}^{1} dx_1 \int_{\tau_m/x_1}^{1} dx_2 [V_i(x_1)V_j(x_2) + (i \leftrightarrow j)] \hat{\sigma}_{V_iV_j\to X}(x_1x_2s) \quad . \tag{3}$$

It is convenient to define the luminosities of bosons inside the fermion as

$$\frac{dL_{ij}}{d\tau} = \int_{\tau}^{1} \frac{dx_1}{x_1} \frac{1}{(1+\delta_{ij})} \left[V_i(x_1)\ V_j(\frac{\tau}{x_1}) + (i \leftrightarrow j)\right] \quad , \tag{4}$$

and Eq. (3) is reduced to

$$\sigma = \int_{\tau_m}^{1} d\tau \frac{dL_{ij}}{d\tau} \hat{\sigma}_{V_iV_j\to X}(\tau s) \quad , \tag{5}$$

where $\tau = \hat{s}/s$, and $\hat{\sigma}_{V_iV_j\to X}$ is the cross section of the subprocess $V_iV_j \to X$.

From (1), (2) and (4) we can easily show that

$$\frac{dL_{\gamma\gamma}}{d\tau} = \left[\frac{\alpha}{2\pi}\ln\left[\frac{\hat{s}}{4m_e^2}\right]\right]^2 \frac{(2+\tau)\ln(\frac{1}{\tau}) - 2(1-\tau)(3+\tau)}{\tau} \quad , \tag{6a}$$

$$\frac{dL_{V_TV_T}}{d\tau} = \left[\frac{\alpha}{2\pi}(C_V^2 + C_A^2)\ln\left[\frac{\hat{s}}{M_W^2}\right]\right]^2 \frac{(2+\tau)^2 \ln(\frac{1}{\tau}) - 2(1-\tau)(3+\tau)}{\tau} \quad , \tag{6b}$$

$$\frac{dL_{V_LV_L}}{d\tau} = \left[\frac{\alpha}{\pi}(C_V^2 + C_A^2)\right]^2 \frac{(1+\tau) \ln(\frac{1}{\tau}) + 2(\tau-1)}{\tau} \quad , \tag{6c}$$

$$\frac{dL_{V_LV_T}}{d\tau} = \left[\frac{\alpha}{2\pi}(C_V^2 + C_A^2)\right]^2 \ln\left[\frac{\hat{s}}{M_T^2}\right] \frac{4(1+\tau) \ln(\frac{1}{\tau}) - (1-\tau)(7+\tau)}{\tau} \quad , \tag{6d}$$

where $M_T = M_V$ for W_T and Z_T, and $M_T = 2m_e$ for the photon.

The above luminosities are basically the quantity of vector bosons that can be found in the electron (or positron), and are plotted in fig. 1. Notice that the electron (and positron) will mostly carry a cloud of transversally polarized bosons, whose luminosity is one order of magnitude larger than the others as can be observed in fig. 1. As we shall verify, this difference may be compensated in the total cross section by the fact that the cross sections for subprocesses involving longitudinally polarized bosons are larger than the ones involving transversally polarized bosons by more than one order of magnitude.

III. CROSS SECTIONS FOR THE ELEMENTARY PROCESS $V_iV_j \to$ LEPTONS

The process we are interested in is $e^+e^- \to e^+e^-(\nu)V_iV_j \to e^+e^-(\nu)X$ where $V_{i(j)}$ can be a photon or a weak boson. In this last case, as we have distribution functions

for longitudinally and transversally polarized bosons, we must compute the cross sections for each one of these polarizations. If the final state (X) is a pair L^+L^-, the initial one (V_iV_j) may be: $\gamma\gamma$, W_TW_T, Z_TZ_T, W_LW_L, Z_LZ_L, γZ_T, γZ_L, W_LW_T and Z_LZ_T (where L(T) means longitudinal (transversal) polarization). When the final state is a lepton $(L^\pm)$ and its neutrino (ν_L), the possible V_iV_j contributions are γW_T, γW_L, W_LZ_T, W_TZ_L. In the following we shall present the main cross sections of the processes quoted above in the limit of high energies; for the complete expressions we refer the reader to refs. (10) and (5).

a) $W_L^+W_L^- \longrightarrow L^+L^-$

$$\hat{\sigma}_{WW\to LL}(\hat{s}) = \frac{\pi\alpha^2}{2\sin^4\theta_W}\left[\frac{M_L}{M_W}\right]^4 \frac{\beta}{\hat{s}}\left\{-1 + \frac{\mathscr{L}}{\beta} + M_H^2(\hat{s}-M_H^2)\,\chi_H\left[-1 + \frac{(1-\beta^2)}{2\beta}\mathscr{L}\right]\right.$$

$$\left. + \frac{\beta^2}{(1-\beta^2)}M_H^4\,\chi_H\right\} \quad , \tag{7}$$

where

$$\beta = \left[1 - \frac{4M_L^2}{\hat{s}}\right]^{1/2} \quad ,$$

$$\mathscr{L} \equiv \ln\frac{(1+\beta)}{(1-\beta)} \quad ,$$

$$\chi_H = \frac{1}{(\hat{s}-M_H^2)^2 + \Gamma_H^2\, M_H^2} \quad .$$

M_L and M_H are respectively the lepton and Higgs boson masses.

b) $Z_L^0 Z_L^0 \longrightarrow L^+ L^-$

$$\underset{ZZ\to LL}{\hat{\sigma}(\hat{s})} = \frac{\pi\ \alpha^2}{2\sin^4\theta_W\ \cos^4\theta_W}\left[\frac{M_L}{M_Z}\right]^4 \frac{\beta}{\hat{s}}\left\{\left[1 + 2M_H^2(\hat{s}-M_H^2)\ \chi_H\right]\left[-1+\frac{1}{2\beta}\mathscr{L}\right]\right.$$

$$\left. + \frac{\beta^2}{(1-\beta^2)} M_H^4\ \chi_H\right\} \quad , \qquad (8)$$

c) $\gamma Z_L^0 \longrightarrow L^+ L^-$

$$\underset{ZZ\to LL}{\hat{\sigma}(\hat{s})} = \frac{\pi\ \alpha^2}{\sin^2\theta_W\ \cos^2\theta_W}(1-4\sin^2\theta_W)^2\left[\frac{M_L}{M_Z}\right]^2 \frac{\beta}{\hat{s}}\left[-2+\frac{1}{\beta}\mathscr{L}\right] \quad . \qquad (9)$$

d) $\gamma W_L^{\mp} \longrightarrow L^{\mp}\ \overset{(-)}{\nu}_L$

$$\underset{\gamma W\to L\nu}{\hat{\sigma}(\hat{s})} = \frac{\pi\ \alpha^2}{2\ \sin^2\theta_W}\left[\frac{M_L}{M_W}\right]^2 \frac{1}{\hat{s}}\left\{-(1-\eta)(1-4\eta) + \left[1-2\eta(1+\eta)\right]\mathscr{L}\right\} \quad , \qquad (10)$$

where

$$\eta = \frac{M_L^2}{\hat{s}} \quad .$$

We would like to recall that the cross sections (7) to (10) are the result of a sum of diagrams, and some of these separately violate unitarity although the complete sum is well behaved at high energies. It is also important to remember that the above cross sections show an enhancement factor $(M_L/M_V)^4$ for those involving the fusion of two longitudinally polarized bosons, or $(M_L/M_V)^2$ when only one longitudinal boson appears. Obviously, for $M_L << M_V$ there is no enhancement, and we expect that processes originated by transversally polarized bosons dominate, since their luminosity in the fermion is larger.

The origin of the enhancement factors $(M_L/M_W)^2$, is due to the high energy behavior of the longitudinal polarization of the weak bosons (ϵ_L^{μ}), whose dominant term is given by k_μ/M_V. When this polarization vector acts on nonconserved axial currents it introduce a factor M_L/M_V in the amplitude for each longitudinal boson.

IV. RESULTS AND CONCLUSIONS

Our numerical results are shown in figs. 2 and 3. In fig. 2 the curve labelled $\gamma\gamma$ is the contribution of the process $e^+e^- \to e^+e^-\gamma\gamma \to e^+e^-L^+L^-$, whose subprocess cross section $\gamma\gamma \to L^+L^-$ is

$$\underset{\gamma\gamma\to LL}{\hat{\sigma}(\hat{s})} = \frac{4\pi\alpha^2}{\hat{s}}\beta\left[\frac{(3-\beta^4)}{2\beta}\ln\left(\frac{1+\beta}{1-\beta}\right) - 2 + \beta^2\right] \quad , \qquad (11)$$

which, among all the process of fusion of vector bosons, is the dominant up to $M_L \approx 300$ GeV. For larger lepton masses$^{(11)}$ the processes involving longitudinally polarized bosons start being more important than the $\gamma\gamma$ one.

Notice that the introduction of at least one longitudinal boson (γW_L), producing a heavy lepton and its neutrino (fig. 2) overcomes the $\gamma\gamma$ production of a pair of heavy leptons. There are two main factors justifying this behavior. Firstly the enhancement $(M_L/M_V)^2$ and secondly the larger phase space for the pair $L\nu$.

The fact that electrons and positrons do not contain too many W's and Z's is also reflected in fig. 3, where we have a enhancement factor $(M_L/M_V)^4$ for $W_L W_L$ and $Z_L Z_L$, but they are not enough to overcome the larger γ luminosity and a smaller enhancement $(M_L/M_V)^2$ of the process γW_L.

In fig. 3 we also notice the effect of Higgs boson exchange diagrams, however larger or smaller masses for the Higgs bosons will not modify the results, because their effects appear in a region where the two–photon process is clearly dominant.

In conclusion, for very heavy leptons, the fusion of a longitudinally polarized charged boson and a photon producing a pair $L\nu$, do win over any other vector boson fusion process, although it is not the main mechanism for heavy lepton production. If such leptons exist they can be better found in the direct annihilation of e^+e^- through the γ and Z^0. The events containing a heavy lepton and its neutrino from vector boson fusion can be detected looking for a jet plus missing energy, although at CLIC energies and for a lepton mass in the range 400 – 500 MeV we shall not have more than $O(10^2)$ events.

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11. There are experimental constraints on the splitting of a new fermion generation, which for a light Higgs boson can be quite stringent, for example, $|M_L - M_\nu| < 310$ GeV for $M_H = 100$ GeV (U. Amaldi et al., *Phys. Rev.* **D36** (1987) 1385), however, these constraints are overcomed in some extensions of the standard model.

FIGURE CAPTIONS

Fig. 1 – **(a)** Luminosities for $\gamma\gamma$, $W_T W_T$ and $Z_T Z_T$ fusion, **(b)** idem for $W_L W_L$ and $Z_L Z_L$, **(c)** idem for γW_T, γZ_T, γW_L and γZ_L, **(d)** idem for $W_L W_T$, $Z_L Z_T$ and $W_L Z_T$.

Fig. 2 – Cross sections of the process $e^+e^- \longrightarrow e^+e^-L^+L^-$ through $\gamma\gamma$ and γZ_L fusion, and for $e^-e^+ \longrightarrow e\,\overset{(-)}{\nu}\,L\,\overset{(-)}{\nu}$ via γW_L fusion.

Fig. 3 – Cross section for the process $e^+e^- \longrightarrow e^+e^-L^+L^-$ via WW and ZZ fusion at $\sqrt{s} = 2$ TeV for $M_H = 1$ TeV (solid curves), $M_H = 500$ GeV (dashed curves), and $M_H = 200$ GeV (dotted curves).

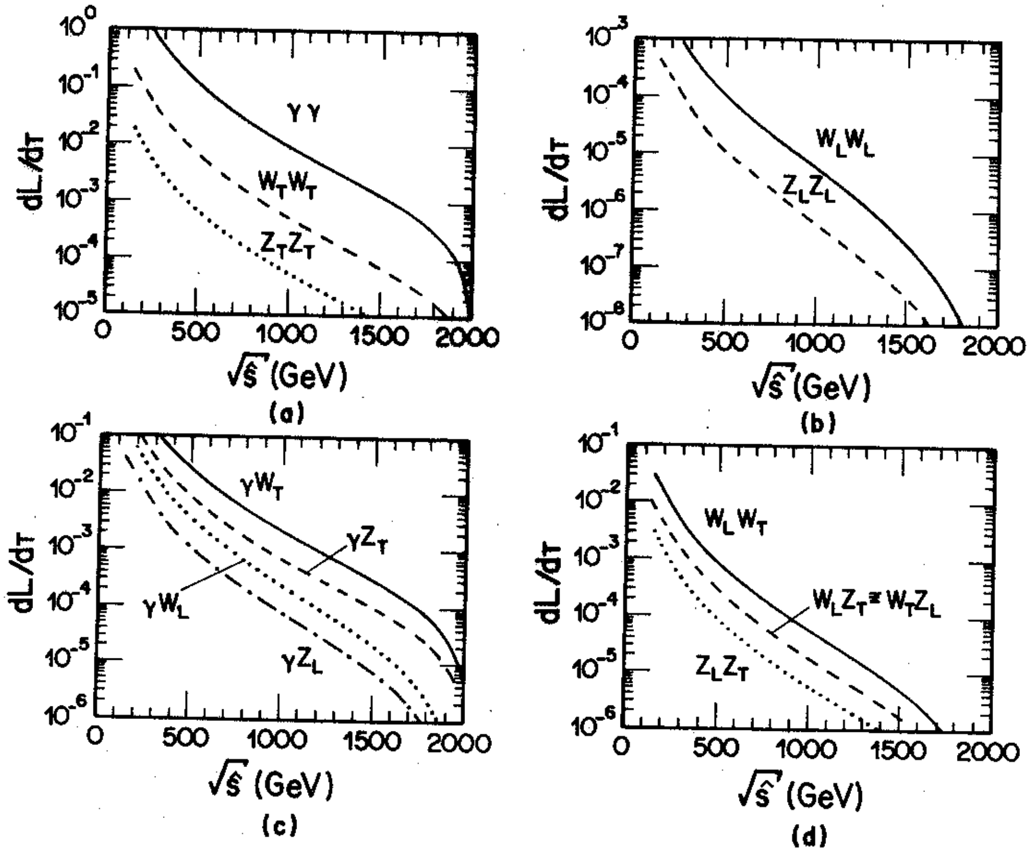
dL/dτ
10^0
10^-1
10^-2
10^-3
10^-4
10^-5
γγ
W_T W_T
Z_T Z_T
0 500 1000 1500 2000
√ŝ (GeV)
(a)
10^-3
10^-4
10^-5
10^-6
10^-7
10^-8
W_L W_L
Z_L Z_L
√ŝ (GeV)
(b)
10^-1
10^-2
10^-3
10^-4
10^-5
10^-6
γW_T
γZ_T
γW_L
γZ_L
√ŝ (GeV)
(c)
W_L W_T
W_L Z_T ≈ W_T Z_L
Z_L Z_T
√ŝ (GeV)
(d)

FIG. 1

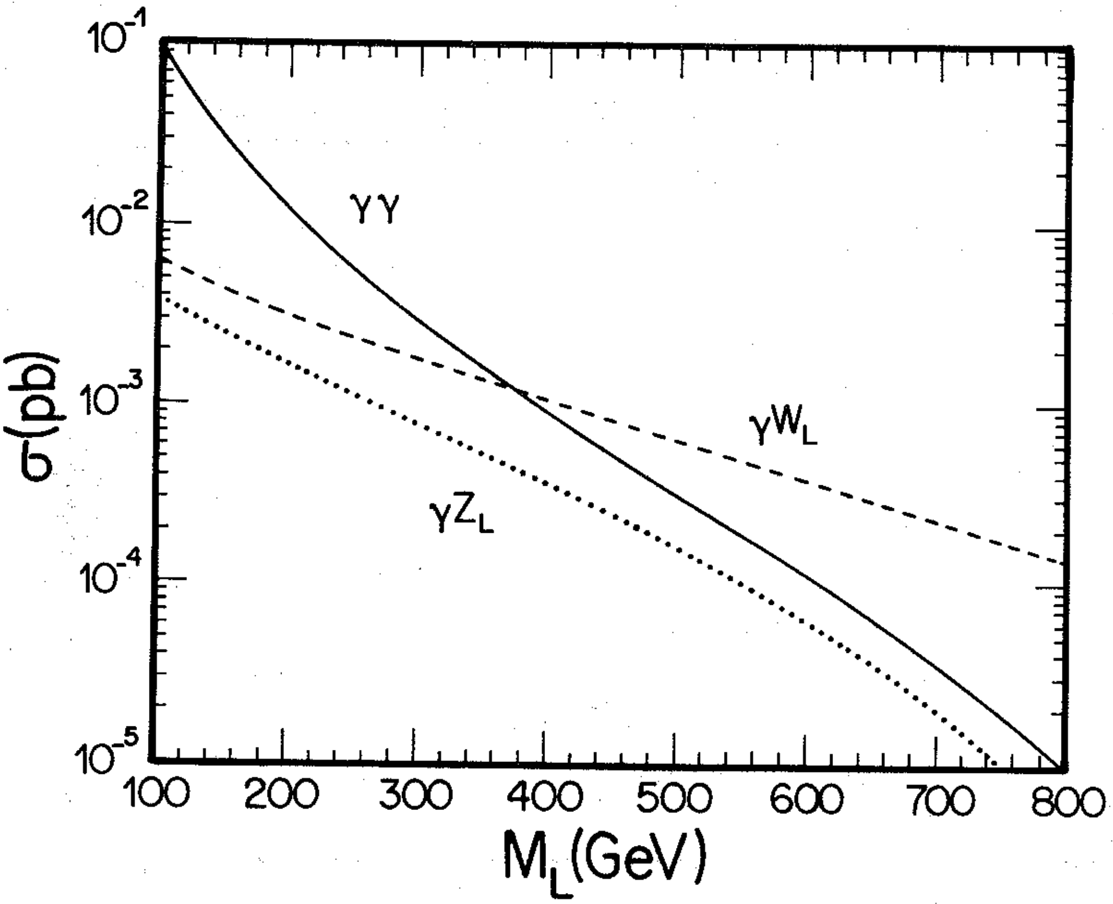
σ(pb)
10^-1
10^-2
10^-3
10^-4
10^-5
γγ
γW_L
γZ_L
100 200 300 400 500 600 700 800
M_L(GeV)

FIG. 2

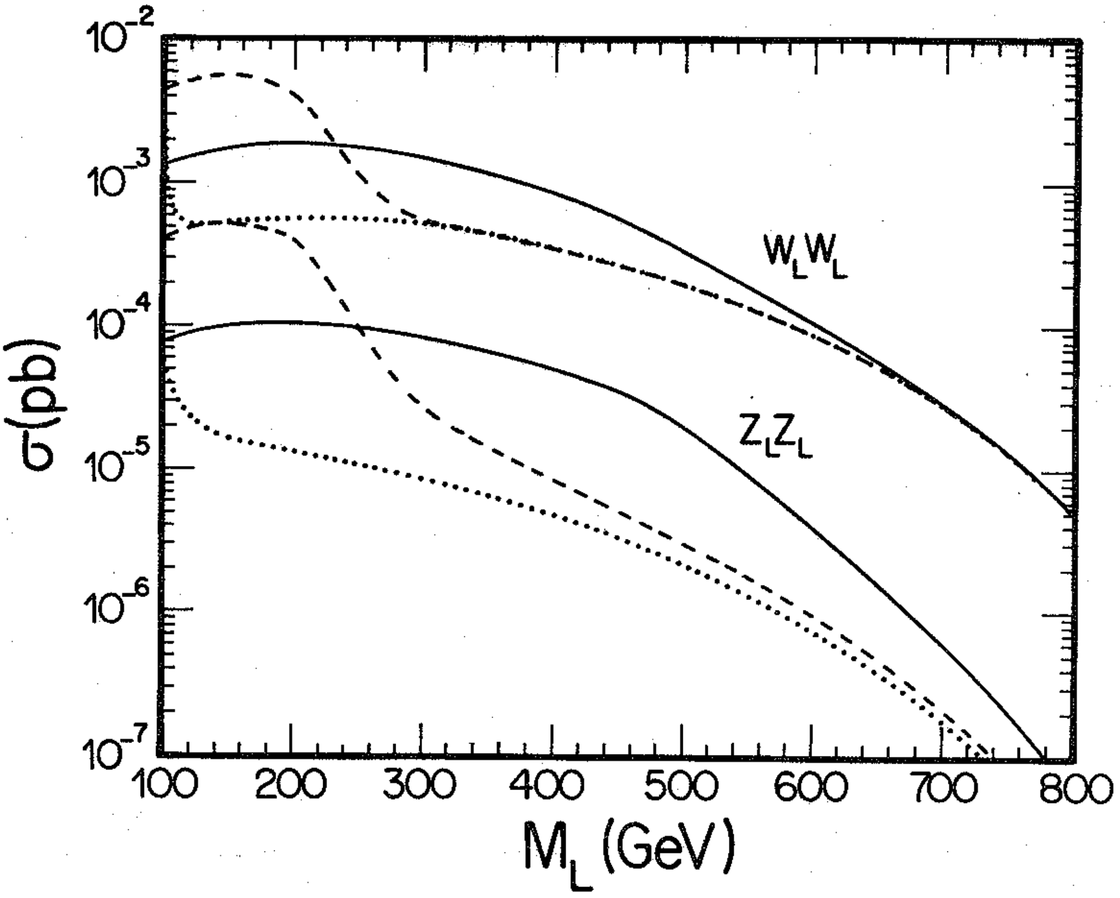
10^{-2}
10^{-3}
10^{-4}
10^{-5}
10^{-6}
10^{-7}
σ(pb)
$W_L W_L$
$Z_L Z_L$
100
200
300
400
500
600
700
800
M_L (GeV)

FIG. 3