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Fermion condensate model of electroweak interactions

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Abstract. A new dynamical symmetry breaking model of electroweak interactions is proposed based on interacting fermions. Two fermions of different $SU_L(2)$ representations form a symmetry breaking condensate and generate the lepton and quark masses. The weak gauge bosons obtain their usual standard model masses from a gauge-invariant Lagrangian of a doublet scalar field composed of the new fermion fields. The new fermion fields become massive by condensation. It is shown that the new charged fermions are produced at the next linear colliders in large number. The model is a low-energy one, which cannot be renormalized perturbatively. For the parameters of the model, unitarity constraints are presented.

One of the most important open questions of particle physics is the origin of electroweak symmetry breaking. Scalars have not been seen in experiments, emphasizing the importance of investigating models beyond the standard one [1–6]. In particular, technicolor models [1] embody dynamical symmetry breaking but it is not easy to fulfill electroweak precision tests. The top quark plays a key role in top condensate models [2], which predict too large masses for the top quark and the Higgs boson. Electroweak symmetry breaking by the condensate of massive matter vector bosons was also put forward [3].

In this note we propose a new low-energy effective model based on four-fermion interaction of new doublet and singlet fermion fields forming condensates. These lead to massive fermions in the linearized approximation and a dynamical symmetry breaking realization is put forward as in the Nambu–Jona Lasinio model [4] and no bare mass terms are introduced in the starting Lagrangian. Non-vanishing lepton and quark masses are coming from gauge-invariant couplings of known and new fermions. The Kobayashi–Maskawa description is unchanged. The usual gauge boson masses follow from the gauge-invariant interaction of a scalar doublet composed of new fermion fields. The production cross section of the new fermions is calculated for the next generation of linear electron–positron colliders. Unitarity constraints for two-particle elastic scattering processes are presented.

We consider the standard model but replace the usual Higgs sector with new fermions having non-renormalizable four-fermion interactions. The new fermions are a neutral singlet,

$$\Psi_S, \qquad T = Y = 0\,,$$

and a weak doublet with hypercharge 1,

$$\Psi_D = \begin{pmatrix} \Psi_D^+ \\ \Psi_D^0 \end{pmatrix}, \qquad T = \frac{1}{2},\, Y = 1\,.$$

Ψ_D^+ $\left(\Psi_D^0\right)$ is a field of positive (zero) electric charge. The doublet has the same quantum numbers as the standard Higgs field. The new fermions are assumed to have four-fermion interactions as a result of an unknown new physics. Let us assume that in the vacuum a symmetry breaking condensate is formed:

$$\left\langle \overline{\Psi}_S \Psi_D \right\rangle_0 = \left\langle \begin{pmatrix} \overline{\Psi}_S \Psi_D^+ \\ \overline{\Psi}_S \Psi_D^0 \end{pmatrix} \right\rangle_0 \neq 0\,. \tag{1}$$

By $SU_L(2)$ transformations we can always transform the upper component into 0. By $U_Y(1)$ transformations of Ψ_D the condensate can be chosen real. $\overline{\Psi}_S \Psi_D$ resembles the standard scalar doublet. The condensate (1) breaks properly the weak $SU_L(2) \times U_Y(1)$ and respects the electromagnetic $U_{\text{em}}(1)$ symmetry. The strength of the condensate is

$$\left\langle \overline{\Psi}_{S\alpha} \Psi_{D\beta}^0 \right\rangle_0 = a_3 \delta_{\alpha\beta}, \qquad \left\langle \overline{\Psi}_{S\alpha} \Psi_{D\beta}^+ \right\rangle_0 = 0\,, \tag{2}$$

where a_3 is real.

In what follows we build up the low-energy effective theory. We replace in the standard model Lagrangian the Higgs sector with gauge-invariant kinetic terms for the new

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fermions and new four-fermion interactions, $L_\Psi + L_f$,

$$L_\Psi = \mathrm{i}\overline{\Psi}_D D_\mu \gamma^\mu \Psi_D + \mathrm{i}\overline{\Psi}_S \partial_\mu \gamma^\mu \Psi_S + \lambda_1 \left(\overline{\Psi}_D \Psi_D\right)^2 + \lambda_2 \left(\overline{\Psi}_S \Psi_S\right)^2 + \lambda_3 \left(\overline{\Psi}_D \Psi_D\right)\left(\overline{\Psi}_S \Psi_S\right) , \quad (3)$$

$$L_f = g_f \left(\overline{\Psi}^f_L \Psi^f_R\right)\left(\overline{\Psi}_S \Psi_D\right) + g_f \left(\overline{\Psi}^f_R \Psi^f_L\right)\left(\overline{\Psi}_D \Psi_S\right) . \quad (4)$$

Here

$$D_\mu = \partial_\mu - \mathrm{i}\frac{g}{2}\underline{\tau}\,\underline{A}_\mu - \mathrm{i}\frac{g'}{2} B_\mu ; \quad (5)$$

$\underline{A}_\mu, B_\mu$ and g, g' are the usual weak gauge boson fields and couplings, respectively.

L_f couples $\Psi_{S,D}$ to the traditional left (right) handed lepton (or quark) doublet Ψ^f_L (singlet Ψ^f_R) and it can be extended to three families. The dimensionful coupling constants can be written as $\lambda_i = \tilde{\lambda}_i/M^2$ and $g_i = \tilde{g}_i/M^2$, where the couplings with tilde are dimensionless and M can be considered as the scale of the new physics. M is expected to be a few TeV. The masses of the weak gauge bosons are coming from effective interactions specified later. λ_i, g_f are assumed to be positive.

The mixed condensate (1) generates masses for the standard quarks and leptons and leads to four-fermion contact interactions. For example, for the electron and electron neutrino doublet it reads

$$L_f = g_e \left(\overline{\nu}\, e_R \overline{\Psi}_S \Psi^+_D + \overline{e}_L e_R \overline{\Psi}_S \Psi^0_D + \overline{e}_R \nu\, \overline{\Psi}^+_D \Psi_S + \overline{e}_R e_L \overline{\Psi}^0_D \Psi_S \right) . \quad (6)$$

In the linearized approximation L_f generates the electron mass

$$m_e = -4 g_e a_3 \quad (7)$$

and $e^+e^-\Psi_S\Psi_D$-type interactions in the physical gauge $\overline{\Psi}_S \Psi^+_D = 0$. Down-type quark masses are generated similarly. For up quarks, as usual, couplings to the charge conjugate $\tilde{\Psi}_D = \mathrm{i}\tau_2 (\Psi_D)^\dagger = \left(\left(\Psi^0_D\right)^\dagger , - \left(\Psi^+_D\right)^\dagger \right)$ must be introduced. Introducing non-diagonal quark bilinears, the Kobayashi–Maskawa mechanism emerges. As in the standard model, from (7) for two particles $m_i/m_j = g_i/g_j$, but a_3 is not determined alone by the Fermi coupling constant.

In order to generate masses for the new fermions, the condensates

$$\left\langle \overline{\Psi}^0_{D\alpha} \Psi^0_{D\beta} \right\rangle_0 = a_1 \delta_{\alpha\beta} , \qquad \left\langle \overline{\Psi}_{S\alpha} \Psi_{S\beta} \right\rangle_0 = a_2 \delta_{\alpha\beta} \quad (8)$$

are introduced, where a_1, a_2 are real. Then, in the linearized approximation one has

$$L_\psi \to L^{\mathrm{lin}}_\Psi = -m_+ \overline{\Psi^+_D} \Psi^+_D - m_1 \overline{\Psi^0_D} \Psi^0_D - m_2 \overline{\Psi}_S \Psi_S - m_3 \left(\overline{\Psi^0}_D \Psi_S + \overline{\Psi}_S \Psi^0_D \right) , \quad (9)$$

with

$$\begin{aligned} m_\pm &= -\left(8\lambda_1 a_1 + 4\lambda_3 a_2\right) = -2a_1\lambda_1 + m_1 , \\ m_1 &= -\left(6\lambda_1 a_1 + 4\lambda_3 a_2\right) , \\ m_2 &= -\left(6\lambda_2 a_2 + 4\lambda_3 a_1\right) , \\ m_3 &= \lambda_3 a_3 . \end{aligned} \quad (10)$$

The condensate of the charged Ψ^+_D would only shift m_+, m_1, m_2 by $-6\lambda_1$, $-8\lambda_1$, $4\lambda_3$ times the condensate. Equations (7) and (10) show that, in general, sign changes of the condensates may influence the sign of the mass parameters. This also happens in [4]. To define physical eigenvalues, (9) is diagonalized by the unitary transformation

$$\begin{aligned} \Psi_1 &= c\,\Psi^0_D + s\,\Psi_S , \\ \Psi_2 &= -s\,\Psi^0_D + c\,\Psi_S , \end{aligned} \quad (11)$$

where $c = \cos\phi$ and $s = \sin\phi$; ϕ is the mixing angle. The masses of the physical fermions Ψ_1, Ψ_2 are

$$2M_{1,2} = m_1 + m_2 \pm \frac{m_1 - m_2}{\cos 2\phi} . \quad (12)$$

The mixing angle is given by

$$2m_3 = (m_1 - m_2)\tan 2\phi . \quad (13)$$

Once $m_1 = M_1$, $m_2 = M_2$, the mixing vanishes, $m_3 = 0$ and vice versa.

It follows that the physical eigenstates themselves form condensates, since

$$\begin{aligned} c^2 \left\langle \overline{\Psi}_{1\alpha} \Psi_{1\beta} \right\rangle_0 + s^2 \left\langle \overline{\Psi}_{2\alpha} \Psi_{2\beta} \right\rangle_0 &= a_1 \delta_{\alpha\beta} , \\ s^2 \left\langle \overline{\Psi}_{1\alpha} \Psi_{1\beta} \right\rangle_0 + c^2 \left\langle \overline{\Psi}_{2\alpha} \Psi_{2\beta} \right\rangle_0 &= a_2 \delta_{\alpha\beta} , \\ cs \left\langle \overline{\Psi}_{1\alpha} \Psi_{1\beta} \right\rangle_0 - cs \left\langle \overline{\Psi}_{2\alpha} \Psi_{2\beta} \right\rangle_0 &= a_3 \delta_{\alpha\beta} . \end{aligned} \quad (14)$$

There is no mixed condensate as Ψ_1, Ψ_2 are independent. For $a_3/cs > 0$ ($M_1 > M_2$), Ψ_1 forms the larger condensate. Combining the equations of (14), one finds

$$a_3 = \frac{1}{2}\tan 2\phi\,(a_1 - a_2) . \quad (15)$$

For $a_1 = a_2$, one cannot put $a_3 \neq 0$ for $\cos 2\phi \neq 0$. As is seen, (15) is equivalent to $\left\langle \overline{\Psi}_{1\alpha} \Psi_{2\beta} \right\rangle_0 = 0$. Comparing (15) to (13) yields

$$m_1 - m_2 = \lambda_3 (a_1 - a_2) . \quad (16)$$

Substituting $m_1 - m_2$ from (10), we are led to the consistency condition

$$a_1 (\lambda_3 - 2\lambda_1) = a_2 (\lambda_3 - 2\lambda_2) . \quad (17)$$

A consequence is that $a_1 \neq a_2$ goes with $\lambda_1 \neq \lambda_2$.

In the physical spectrum there are three new fermions, a charged one and two neutral ones. They have many self-interactions given by (3) and (11), like $\Psi^+_D \Psi^+_D \Psi^-_D \Psi^-_D$, $\Psi^+_D \Psi^-_D$ ($\Psi_1\Psi_1$, $\Psi_2\Psi_2$, $\Psi_1\Psi_2$), Ψ_1^4, $\Psi_1^3\Psi_2$, $\Psi_1^2\Psi_2^2$, $\Psi_1\Psi_2^3$, Ψ_2^4.

The masses of the weak gauge bosons arise from the effective interactions of an auxiliary composite $Y=1$ scalar doublet,

$$\Phi=\begin{pmatrix}\Phi^{+}\\ \Phi^{0}\end{pmatrix}=\overline{\Psi}_S\Psi_D\,. \tag{18}$$

Φ develops a gauge-invariant kinetic term in the low-energy effective description

$$L_H=h\left(D_\mu\Phi\right)^{\dagger}\left(D^\mu\Phi\right)\,. \tag{19}$$

Here D_μ is the usual covariant derivative (5).

The coupling constant h sets the dimension of L_H, $[h]=-4$ in mass dimension. We assume that $h>0$. Equation (19) is a non-renormalizable Lagrangian leading to the weak gauge boson masses and some of the interactions of the new fermions with the standard gauge bosons.

In the gauge $\Phi^{+}=0$, L_H can be written as

$$\begin{aligned}h^{-1}L_H=&\frac{g^2}{2}W^{-}_{\mu}W^{+\mu}\Phi^{0\dagger}\Phi^{0}+\frac{g^2}{4\cos^2\theta_W}Z_\mu Z^\mu\Phi^{0\dagger}\Phi^{0}\\&+\Bigg[\partial^\mu\Phi^{0\dagger}\partial_\mu\Phi^{0}-\frac{\mathrm{i}}{2}\frac{g}{\cos\theta_W}\left(\partial^\mu\Phi^{0\dagger}\right)\Phi^{0}Z_\mu\\&\qquad+\frac{\mathrm{i}}{2}\frac{g}{\cos\theta_W}\Phi^{0\dagger}Z_\mu\left(\partial^\mu\Phi^{0}\right)\Bigg]\end{aligned} \tag{20}$$

in terms of the usual vector boson fields.

In the linearized approximation of (20), we put

$$h\,\Phi^{0\dagger}\Phi^{0}\to h\left\langle\Phi^{0\dagger}\Phi^{0}\right\rangle_0=h\left(16a_3^2-4a_1a_2\right)=\frac{v^2}{2}\,, \tag{21}$$

leading to the standard masses

$$m_W=\frac{gv}{2}\,,\qquad m_Z=\frac{gv}{2\cos\theta_W}\,. \tag{22}$$

v^2 is, as usual, $\left(\sqrt{2}G_F\right)^{-1}$; $v=254\,\mathrm{GeV}$. The tree masses naturally fulfill the important relation $\rho_{\mathrm{tree}}=1$. Equations (21) and (15) impose the restriction

$$v^2=8h\left[\tan 2\Phi\left(a_1-a_2\right)^2-a_1a_2\right]>0\,. \tag{23}$$

Equation (20) describes several non-renormalizable interactions of the type Z, Z^2, W^2 times Ψ_S^2, $\left(\Psi_D^0\right)^2$ or $\Psi_S^2\left(\Psi_D^0\right)^2$. Their strengths in h units are determined by the gauge coupling constants.

A transparent experimental consequence of the model is provided by the doublet kinetic term in (3), whence we obtain the usual renormalizable fermion pair couplings to gauge bosons:

$$\begin{aligned}L^I=&\overline{\Psi_D^{+}}\gamma^\mu\Psi_D^{+}\left(eA_\mu-e\cot 2\theta_W Z_\mu\right)\\&+\frac{g}{2\cos\theta_W}\overline{\Psi_D^{0}}\gamma^\mu\Psi_D^{0}Z_\mu\\&+\frac{g}{\sqrt{2}}\left(\overline{\Psi_D^{+}}\gamma^\mu\Psi_D^{0}W_\mu^{+}+\overline{\Psi_D^{0}}\gamma^\mu\Psi_D^{+}W_\mu^{-}\right)\,.\end{aligned} \tag{24}$$

To obtain the physical interactions for the neutral component, the mixing (11) must be taken into account; then

$$\begin{aligned}L^I=&\overline{\Psi_D^{+}}\gamma^\mu\Psi_D^{+}\left(eA_\mu-e\cot 2\theta_W Z_\mu\right)\\&+\frac{e}{\sin 2\theta_W}Z_\mu\Big(c^2\overline{\Psi}_1\gamma^\mu\Psi_1+s^2\overline{\Psi}_2\gamma^\mu\Psi_2\\&\qquad-sc\left(\overline{\Psi}_1\gamma^\mu\Psi_2+\overline{\Psi}_2\gamma^\mu\Psi_1\right)\Big)\\&+\left[\frac{g}{\sqrt{2}}W_\mu^{+}\left(c\overline{\Psi_D^{+}}\gamma^\mu\Psi_1-s\overline{\Psi_D^{+}}\gamma^\mu\Psi_2\right)+h.c.\right]\,.\end{aligned} \tag{25}$$

The interaction L_f in (6) turns out to be very weak. Indeed, from (7) and (21) we have an upper bound for g_e, $g_e\le\sqrt{2h}\frac{m_e}{v}=\sqrt{2h}g_e^{\mathrm{SM}}$, which is suppressed by two factors of the scale of new physics compared to the standard model value g_e^{SM}.

To test the model at the forthcoming accelerators, we consider the production of new fermion pairs in electron–positron annihilation. It is most useful to investigate the case of a charged new fermion pair; we denote this by $D^{+}D^{-}$.

The contact graph from (6) yields the cross section

$$\sigma\left(e^{+}e^{-}\to D^{+}D^{-}\right)=\frac{g_e^2}{16\pi}s\sqrt{1-4\frac{m_{+}^2}{s}}\left(1-\frac{5}{2}\frac{m_{+}^2}{s}\right)\,, \tag{26}$$

where s is the center of mass energy squared. The cross section is negligible at moderate s. For example, at $h\sim(2\,\mathrm{TeV})^{-4}$ and $\sqrt{s}=1\,\mathrm{TeV}$ it is still in the order of 10^{-13} fb.

We expect a higher number of events from the photon and Z exchange processes $e^{+}e^{-}\to\gamma,Z\to D^{+}D^{-}$. The usual SM coupling at the $e^{+}e^{-}Z$ vertex is

$$\mathrm{i}\frac{g}{2\cos\theta_W}\gamma_\mu\left(g_V+\gamma_5 g_A\right)\,,$$

$$g_V=-\frac{1}{2}+2\sin^2\theta_W\,,\quad g_A=-\frac{1}{2}\,.$$

By making use of (25), one obtains the cross section

$$\sigma\left(e^{+}e^{-}\to D^{+}D^{-}\right)=\frac{1}{16\pi}\sqrt{1-4\frac{m_{+}^2}{s}}\frac{1}{s}\left|M\right|^2\,,$$

$$\begin{aligned}\left|M\right|^2=&\tfrac{4}{3}e^4\tfrac{s+2m_{+}^2}{s}+\tfrac{2}{3}\tfrac{e^4}{\sin^2\theta_W\cos^2\theta_W}g_V\tfrac{s+2m_{+}^2}{s-m_Z^2}\\&+\tfrac{1}{12}\tfrac{e^4}{\sin^4\theta_W\cos^4\theta_W}\left(g_V^2+g_A^2\right)s\tfrac{s+2m_{+}^2}{\left(s-m_Z^2\right)^2}\,,\end{aligned} \tag{27}$$

where the three terms in $|M|^2$ are coming from photon exchange, photon–Z interference and pure-Z exchange. A similar cross section belongs to the neutral pair productions, too. The cross section rises fast after the threshold; at high energies it falls off as $1/s$, reflecting that all the interactions are renormalizable in the process. The cross section is given in Table 1 for a few masses and plotted vs.

Table 1. Cross section of D^+D^- production at $\sqrt{s}=500$ GeV

m_+ (GeV)	100	150	200
$\sigma\left(e^+e^-\to D^+D^-\right)$ (fb)	560	535	450

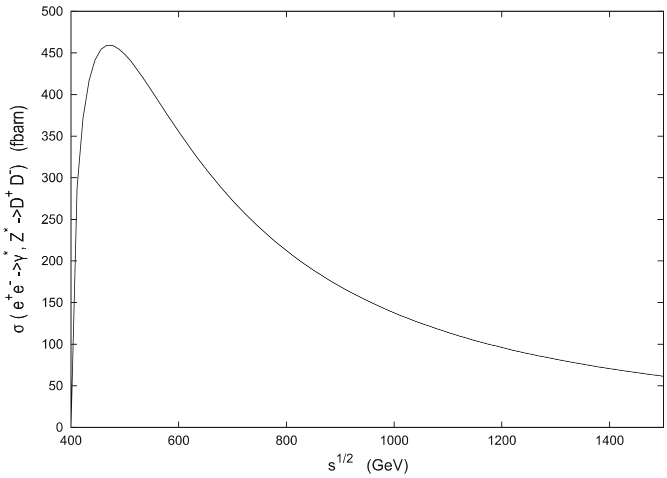

Fig. 1. Cross section of D^+D^- production at electron–positron collider vs. $\sqrt{s}$ for $m_+=200$ GeV

$\sqrt{s}$ in Fig. 1. At a linear collider of $\sqrt{s}=500$ GeV (TESLA) and integrated luminosity 50 fb^{-1}/year a large number of events is expected.

The cross section at $\sqrt{s}=1500$ GeV is an order of magnitude smaller, but with an integrated luminosity of 100 fb^{-1} per annum a large number of events appears and a higher mass range can be searched for, see Table 2.

Now, we provide a rough estimate of the model parameters by imposing partial-wave unitarity, $|\mathrm{Re}\, a_0|\le\frac{1}{2}$, for the $J=0$ partial-wave amplitudes of $2\to2$ processes [7, 8]. Consider D^+D^- elastic scattering. At high energies, $s\gg 4m_+^2$, the contact graph is expected to give the dominant part of a_0, leading to $\lambda_1 s\le 8\pi$. At a maximum possible energy of $s=(1\text{–}25)$ TeV^2, $\lambda_1\le(2.5\text{–}0.1)G_F$. In the case of no mixing we expect a similar upper bound for λ_2, too. To obtain an upper bound for h, let us consider the process $\overline{\Psi}_1\Psi_2\to ZZ$ with longitudinally polarized Z, and $(+-)$ helicity for the incoming fermions. The relevant contact interaction follows from (19):

$$L_H=-Z_\mu Z^\mu\overline{\Psi^0_D}\Psi_S\frac{hg^2a_3}{\cos^2\theta_W}+h.c.\,, \tag{28}$$

where $16a_3^2h\gtrsim\frac{v^2}{2}$. Choosing $\overline{\Psi}_1\Psi_2$ at asymptotic s the unitarity imposes $c^2\left(s^3h\right)^{1/2}\le16\pi v$. This shows that h is in general a very weak coupling: at the scale $s=(1\text{–}25)$ TeV^2, $c^4h\lesssim2\left(1-5^{-6}\right)G_F^2$. We checked that $\bar{t}t\to\overline{\Psi}_1\Psi_2$ provides a weaker bound. Fixing λ_1, λ_2, larger a_1, a_2 give larger masses m_+, m_1, m_2. For instance, at $\lambda_3=0$, $m_+=100\,[300]$ GeV needs a condensate a_1,

$$|a_1|\ge\left(\frac{1}{60}-\frac{1}{2.4}\right)G_F^{-3/2}\left[\left(\frac{1}{20}-\frac{5}{4}\right)G_F^{-3/2}\right]$$

for $\lambda_1\le(2.5\text{–}0.1)G_F$. The lower bound for $|a_2|$ lies not very far from that of a_1, for m_1 not very far from m_2. It also follows that the leptonic coupling constant g_e is less than its standard model value, $g_e^{\mathrm{SM}}\times\mathcal{O}(G_F)$, if c is not very small. Finally, since $8\,|a_3|\ge2^{1/4}\left(G_Fh\right)^{-1/2}$, one obtains for $s=(1\text{–}25)$TeV^2 $|a_3|\gtrsim\frac{1}{\sqrt{2}}(1\text{–}125)G_F^{-3/2}$ at small mixing.

Similarly to [4], the gap equations impose a lower bound on the coupling constants λ_i. Consider first the $m_3=\lambda_3a_3$ equation in (10). Expressing m_3 with M_1-M_2 by (12) and (13) and a_3 by (14), and making use of the free-field approximation for the condensates of Ψ_1,Ψ_2, we obtain

$$M_1-M_2=\lambda_3\left(I_1-I_2\right)\,, \tag{29}$$

where I_i means the familiar momentum integral

$$I_i=-\frac{M_i}{2\pi^2}\left(\Lambda^2-M_i^2\ln\left(1+\frac{\Lambda^2}{M_i^2}\right)\right)\,, \tag{30}$$

with a cutoff Λ. It is easy to see that for $M_1=M_2$ no restriction occurs for λ_3, but for $M_1\ne M_2$ one obtains $|\lambda_3|\ge2\pi^2\Lambda^{-2}$. Furthermore, around the critical minimum $|\lambda_3|$, $M_{1,2}\ll\Lambda$ and increasing $|\lambda_3|$ increases the masses (as in [4]).

The other equations in (10) can clearly be translated into gap equations expressing $m_{1,2}$ by $M_{1,2}$ and a_i by free-field functions (30). $M_1=0$ and $M_2=0$ are always a solution. The non-zero mass solution of the gap equations imposes $\lambda_{1,2}\ge\pi^2/3\Lambda^2$. λ_3 close to its critical value and both λ_1 and λ_2 near to $5\pi^2/3\Lambda^2$ provide small masses $M_{1,2}\ll\Lambda$ and higher masses attract higher λ_i. Together with the unitarity constraint, one has $\lambda_{1,2}\in[\pi^2/3,8\pi]\Lambda^{-2}$.

In conclusion, we proposed a new dynamical symmetry breaking of the electroweak symmetry based on four-fermion interactions of hypothetical doublet and singlet fermions. Generating masses for all the standard and new particles follows by condensation. Condensates are free parameters of the model. The interactions of new fermions tend to be weak, since the model is a non-renormalizable effective low energy one. From the experimental point of view, $e^+e^-\to D^+D^-$ is advantageous, following from purely gauge interactions, and it results in a large number of events at the next generation of linear colliders.

Table 2. Cross section of D^+D^- production at $\sqrt{s}=1500$ GeV

m_+ (GeV)	100	200	400	700
$\sigma\left(e^+e^-\to D^+D^-\right)$ (fb)	62	61	60	32

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