Electroweak Lepton-Lepton and Lepton-Antilepton Bound States

R.A. Alanakyan

*Theoretical Physics Department, Yerevan Physics Institute, Alikhanian Brothers St.2,*

*Yerevan 375036, Armenia*

E-mail: alanak @ lx2.yerphi.am

**Abstract**

In model independent way we consider the possibility of the existence of fermion-antifermion, fermion-fermion bound states which appear due to $\gamma, Z^0(W^\pm)$-bosons and scalar, pseudoscalars exchanges including radiative corrections. Our consideration includes the case where at least one particle in the bound state is Majorana fermion or scalar. This paper considers various types of bound states of particles contained in various extensions of the Glashow-Weinberg -Salam model. We calculate long-range forces induced by photonic, gluonic and fermionic loop corrections in scalar (pseudoscalar) exchange.
1. Introduction

In this article by model independent way we consider possibility of existence of fermion-antifermion, fermion-fermion bound states which appear due to $\gamma, Z^0(W^\pm)$-bosons and scalar, pseudoscalars exchanges including radiative corrections (long-range potential induced by neutrinos loop (i.e. $Z(W)$-exchanges including box diagrams) has been considered in [1]-[6], long-range forces mediated by Higgs and Goldstown bosons exchange [7] (via loop of very light pseudoscalars):

We also consider (sect.4) a case where at least one of the particles is Majorana neutrino (or neutralino) ($l_i^\pm \bar{\nu}_j, l_i^\pm \bar{N}_j, \nu_i \bar{\nu}_j, \nu_i N_j,, \chi_i^0 \chi_j^0, \tilde{\chi}_i^0 \chi_j^\pm$ etc bound states) which contains in many extensions beyond the Standard Model e.g. in the Minimal Supersymmetric Standard Model [8],[9], or in Left-Right Symmetric Model [10]-[16].

Our consideration ia also true for bound state of scalar neutrinos \footnote{The bound state $\tilde{\nu}\tilde{\nu}^*$ has been considered in [17], however in this paper the authors do not considered $Z^0$-exchanges and also do not consider radiative corrections to Higgs exchange which as we will see below radically change the Higgs potentials.}

As known, at tree level level $Z^0$-bosons exchange lead to Yukawa potential :

$$Z_0(r) = \frac{eg_V^e \exp(-m_Zr)}{4\pi r} \tag{1}$$

where $g_V^e = \frac{1}{s_{WcW}}(T_c - 2Qcs_{W}^2)$ (if attractive center is fermion), $s_W, c_W$-cosinus and sinus of Weinbergs angle, $T_c, Q_c$-are isospin and charge of attractive center. In [1] it has been shown that at large distances the potential enhanced due to neutrino -pairs exchanges:

$$V(r) = \frac{G_F^2}{4\pi^2 r^5} \tag{2}$$
Thus, radiative corrections makes the potential deeper and wider and chances that bound state may exist are increasing. We extend in this paper also the result [1] in range \( r \sim \frac{1}{m_Z} \) or smaller \( r \).

As we will see below analogous behaviour takes place also for Higgs potential: at small distances it is Yukawa-like, at large distances radiative corrections lead to slower decreasing \( \sim r^{-5} \), and to behaviour \( \sim r^{-7}, r^{-9} \) for scalar (pseudoscalar) exchange for correction via photonic and fermionic loop.

### 2. Potential

As mentioned above at large distances \( r >> \frac{1}{m_Z} \) for potential between two electrons in accordance with [1] is described by (2).

At distances \( r < \frac{1}{m_\nu} \) the contribution of neutrinos with masses \( m_\nu \) is suppressed as \( \sim \exp(-2m_\nu r) \).

If we take into account not only neutrino [1] but all light fermions in loop we obtain at \( \frac{1}{2m_f} << r << \frac{1}{m_Z} \):

\[
Z_0(r) = \frac{3e g_\nu^2}{\pi^2 r^5} \frac{\Gamma}{m_Z^3}
\]  

(3)

where \( \Gamma \) is the width of \( Z^0 \)-boson. If in this formula we take into account only neutrinos in loop and besides consider box and triangle diagramm we turn to formula (2).

At small distances \( r << \frac{1}{m_Z} \) from (5) as expected may be obtained well known result:

\[
V(r) = \frac{e g_\nu^2}{4\pi r}(1 + \frac{2}{\pi m_Z} \log(rm_Z))
\]  

(4)

In fact, in (3),(4) all light fermions contributions are included. It must be noted also that in formulas (3),(4) we take into account only \( t \)-channel diagram (Fig.1) in which all \( N \approx \Gamma/\Gamma(Z^0 \rightarrow \nu \nu) \) flavours contribute, whereas
only one flavour contributes into the box and triangle diagramm. Both formulas (3), (4) may be obtained from the general expression (in contrast to the [1] we do not neglect $t'$ in $Z^0$-bosons propagators):

$$Z_0(r) = -\frac{eg_\gamma}{4\pi^2r} \int_0^\infty dx \frac{ImP_T(t) \exp(-r\sqrt{t})}{(t - m_Z^2)^2 + (ImP_T(t))^2},$$

(5)

where at $t >> 4m_f^2$ for imaginary part of polarisation transversal operator we have:

$$ImP_T = -\frac{\Gamma}{m_Z} t$$

(6)

For instance, if $r \gg \frac{1}{m_Z}$ for $t'$s in propagators may be neglected and we obtain (3). At $r \gg \frac{1}{m_Z}$ the main contribution in integral (5) comes from the range where $r\sqrt{t} \ll 1$. Besides, at small distances we can neglect light fermion masses $4m_f^2$ and integrate in (5) within the limits $0 < t < r^{-2}$ and put $e^{-r\sqrt{t}} \approx 1$. After the integration we obtain the result (4).

At large distances $r \sim \frac{1}{m_f}$ we must take into account fermions masses in loop, however, as in previous reference, again $q^2$ and $P_T$ in propagator may be neglected, and analogously [5] we obtain

$$V(r) = \frac{3ag_\gamma}{2\pi^3r^2m_Z^2} \sum_f \frac{\Gamma(Z \rightarrow f\bar{f})}{\Gamma} K_2(2m_f r),$$

(7)

where $K_2(x)$ is modified Bessel function. At distances $r > \frac{1}{m_f}$ the contribution of the flavour $f$ falls as $\sim e^{-2rm_f}$.

Using the general formula (5) we can obtain also long-range forces mediated by virtual photons (virtual gluons) loops [8]:

$$H^0(P^0) \rightarrow \gamma^*\gamma^* \rightarrow H^0(P^0),$$

$$H^0(P^0) \rightarrow gluon^*gluon^* \rightarrow H^0(P^0),$$

or by light fermions loop:
\( H^0(P^0) \rightarrow \bar{f}^* f^* \rightarrow H^0(P^0) \).

This exchanges are shown on Fig.3

Substituting imaginary part of \( H^0(P^0) \)-bosons polarisation operator:

\[
Im P_{H,P} = -\frac{\Gamma(H^0(P^0) \rightarrow \gamma\gamma\gamma)}{m_{H(P)}^3} t^2
\]

into (5) we obtain (for Higgs bosons decays see e.g. [19] and references therein, for \( \pi^0 \) decays see [20], [21]):

\[
V_{S,P}(r) = -a_{S(P)}(1)a_{S,P}(2) \frac{5!}{2\pi^2} \frac{\Gamma(H^0(P^0) \rightarrow \gamma\gamma\gamma)}{(m_{H(P)}r)^7}
\]

Thus, instead of Yucawa potential from tree exchange we have long-range forces which decrease as \( \sim r^{-7} \). For arbitrary \( r \) take place formula like (5). The potentials \( V_{S,P} \) are contains into relativistic equations (see below).

In case of \( \pi^0 \)-meson (which is in fact a composite pseudoscalar) formula true only at \( r > m_{\pi}^{-1} \), at \( r < m_{\pi}^{-1} \) the compositness of \( \pi^0 \)-meson became considerable.

In non-relativistic approximation for \( P^0 \)-bosons exchanges between fermions we have:

\[
V_P(r) = -a_P(1)a_P(2) \frac{5!}{2\pi^2} \frac{1}{4m_1m_2} \frac{\Gamma(P^0 \rightarrow \gamma\gamma\gamma)}{(m_{P^0}r)^9} (-7\vec{\sigma_1}\vec{\sigma_2} + \frac{63(\vec{\sigma_1}\vec{r})(\vec{\sigma_2}\vec{r})}{r^2})
\]

where \( \sigma_{1,2}, m_{1,2} \) are operators of spins and masses of the fermions, \( a_{S,P}(1), a_{S,P}(2) \) its couplings with \( P^0 \)-bosons.

Thus, we have long range potential which fall as \( r^{-9} \).

For fermion loop-induced Higgs potentials we obtain the following model-independent result (at \( r >> \frac{1}{2m_f} \)):

\[
V_{S,P}(r) = -\frac{3a_{S,P}(1)a_{S,P}(2)\Gamma(H^0(P^0) \rightarrow f\bar{f})}{\pi^2 m_{H(P)}^{35}}
\]
Again, at \( r < \frac{1}{2m_f} \), the potentials (12) suppressed as \( \sim e^{-2m_f r} \).

In this form both potentials induced by fermionic loops are contains in equations (19)-(25) below.

In non-relativistic approximation however pseudoscalar exchange (with fermios loop) between fermions has the form:

\[
V_P(r) = -a_P(1)a_P(2) \frac{5!}{2\pi^2} \frac{1}{4m_1m_2} \frac{\Gamma(P^0 \rightarrow f\bar{f})}{(m_P r^2)} \left( -5\vec{\sigma}_1 \vec{\sigma}_2 + \frac{35(\vec{\sigma}_1 \vec{r})(\vec{\sigma}_2 \vec{r})}{r^2} \right)
\]

(12)

while the Higgs potential in non-relativistic approximation again defined by formula (9).

In [18] is given comparison of long-range forces (10) induced by \( \pi^0 \)-meson exchange with various other kinds of long-range forces including Van-Der-Vaals potential.

The presented above consideration for scalar and pseudoscalars exchanges is model independent and applicable also to potentials created by scalars leptons exchanges ((s)fermion -(s)fermion-scalar lepton vertexes arising in theories with \( R \)-parity violation (see ref.[22]-[28]). In this case \( a_{S,P} \) may be not proportional to fermion masses.

3. Dirac Particle Case

The Dirac equation\(^2\) for particle in spherically symmetric potentials \( Z_0, A_0, V_{S,P} \) has the form:

\[
(\hat{k} - e\hat{Z}_0(g_V + g_A\gamma_5) - eQ\hat{A} - m + a_S V_S + a_P \gamma_5 fV_P)u(k) = 0,
\]

(13)

Here \( Z_0(r), A_0(r), V_{S,P}(r) \) are potentials created by immovable center and described in sect.2, \( g_V = \frac{1}{c_W s_W}(T - 2Q s_W^2), g_A = \frac{1}{c_W s_W}2T \), where \( T, Q \) are

\(^2\)for Dirac equation in \( Z^0 \)-boson field for Dirac and Majorana spinors see e.g. [29] and references therein.
isospin and charge of particle respectively.

In (13) $V_{S,P}$ defined in sect.2 above, $a_{S,P}$ are model independent Yukawa couplings of scalars and pseudoscalars with fermions. In $a_{S,P}$ we include also scalars (pseudoscalars) interaction with attractive center.

If we consider a moving antiparticle instead of a particle we must make the replacements in above Dirac equation:

$$g_V, g_A, Q \rightarrow -g_V, -g_A, -Q$$

(14)

Let us consider the movement of fermion (antifermion) in spherically symmetric potentials $A_0, Z_0, V_{S,P}$ using the method developed in[30][31](see also references therein).

Because electroweak interactions violated $P$-parity (term $\gamma_\mu \gamma_5$ in Dirac equation), spinors $\phi, \chi$ must be expressed through linear combination of spheric spinors $\Omega_{j,M}(\vec{n}), \Omega_{j'M}(\vec{n})$ which have different $P$-parity:

$$\psi^T = (\phi, \chi),$$

$$\phi = f_1(r)\Omega_{j,M}(\vec{n}) + (-1)^{l + \frac{1}{2}} f_2(r)\Omega_{j'M}(\vec{n}),$$

$$\chi = g_1(r)\Omega_{j,M}(\vec{n}) + (-1)^{l + \frac{1}{2}} g_2(r)\Omega_{j'M}(\vec{n})$$

(15)

where $l = j \pm \frac{1}{2}$, $l' = 2j - l$. Using identities:

$$\vec{\sigma}\vec{p}\phi = (g'_1(r) + \frac{1 + \kappa}{r} g_1(r))\Omega_{j'M}(\vec{n}) - (g'_2(r) + \frac{1 - \kappa}{r} g_2(r))\Omega_{jM}(\vec{n})$$

(16)

$$\vec{\sigma}\vec{p}\chi = (f'_1(r) + \frac{1 + \kappa}{r} f_1(r))\Omega_{j'M}(\vec{n}) - (f'_2(r) + \frac{1 - \kappa}{r} f_2(r))\Omega_{jM}(\vec{n}),$$

(17)

where

$$\kappa = l(l + 1) - j(j + 1) - \frac{1}{4},$$

(18)

for radial functions we obtain:

$$(g'_2(r) + \frac{1 - \kappa}{r} g_2(r)) + (E - M(r) - V(r))f_1(r) + V_-(r)g_1(r) = 0$$

(19)
\begin{align}
(g'_1(r) + \frac{1 + \kappa}{r} g_1(r)) - (E - M(r) - V(r)) f_2(r) + V_-(r) g_2(r) &= 0 \quad (20) \\
(f'_2(r) + \frac{1 - \kappa}{r} f_2(r)) + (E + M(r) - V(r)) g_1(r) + V_+(r) f_1(r) &= 0 \quad (21) \\
(f'_1(r) + \frac{1 + \kappa}{r} f_1(r)) - (E + M(r) - V(r)) g_2(r) - V_+(r) f_2(r) &= 0 \quad (22)
\end{align}

\begin{equation}
M(r) = m - a_S V_S(r) \quad (23)
\end{equation}

\begin{align}
V(r) &= e g_V Z_0(r) + e Q A_0(r), \quad (24) \\
V_\pm(r) &= e g_A Z_0(r) \pm a_P V_P(r), \quad (25)
\end{align}

For large $r$ we can drop all fields and obtain divergent spheric waves at infinity.

Although above we consider stationar solutions, our consideration is also applicable for time-dependent solutions. In this case we must in (19)-(22) make the replacements:

\begin{equation}
E \to -i \frac{d}{dt}, \quad (26)
\end{equation}

and

\begin{align}
f_i(r) &\to f_i(r, t), \quad (27) \\
g_i(r) &\to g_i(r, t), \quad (28)
\end{align}

At small $r \ll \frac{1}{m\zeta}$ we can put $Z_0(r) = \frac{e g_V}{4\pi r^2}, V_{S,P}(r) = \frac{1}{r}, E = 0$. For radial wave functions we suppose that the solutions have the following form:

Substituting (29) into (19)-(25) we obtain:

\begin{align}
g_i(r) &= a_i r^\gamma, f_i(r) = b_i r^\gamma, \quad (29) \\
(\gamma + 1 - \kappa) a_2 + (a_S - \alpha (g_V g_V^c + Q Q c)) b_1 + a_1 (\alpha g_A g_V) - a_P) &= 0 \quad (30) \\
(\gamma + 1 + \kappa) a_2 - (a_S - \alpha (g_V g_V^c + Q Q c)) b_2 + a_2 (\alpha g_A g_V^c) - a_P) &= 0 \quad (31)
\end{align}
\[(\gamma + 1 - \kappa)b_2 + (-a_S - \alpha(g_V g_V^c + QQ)\alpha_1 + b_1(\alpha g_A g_V^c + a_P) = 0 \quad (32)\]
\[(\gamma + 1 + \kappa)b_1 + (a_S + \alpha(g_V g_V^c + QQ)\alpha_2 - b_2(\alpha g_A g_V^c + a_P) = 0 \quad (33)\]

The determinant of this system of equations must be equal zero, and that condition gives us the equation for defining \(\gamma\).

The general solution of this equation is very complicated. For simplicity consider several cases:
1)\(a_{S,P} = Q = 0\).
2) Only \(a_S\) is nonzero.
3) only \(a_P\) is nonzero.

In case of 1) we obtain 4 solutions:
\[
s = -1 + \sqrt{\kappa^2 - \alpha^2(g_V^2 \pm g_A^2)} \quad (34)
\]
\[
s = -1 - \sqrt{\kappa^2 - \alpha^2(g_V^2 \pm g_A^2)} \quad (35)
\]
In case of 2), 3) the full system of equations decouples on two subsystems of equations with solutions:
\[
s = -1 \pm \sqrt{\kappa^2 - a_{S,P}^2} \quad (36)
\]
\[
s = -1 \pm \sqrt{\kappa^2 + a_{S,P}^2} \quad (37)
\]
In case of pure QED we obtain again \(s = -1 \pm \sqrt{\kappa^2 - Q^2 Q_e^2 \alpha^2}\).

Obviously, if \(\alpha \ll 1\), the solutions (35) must be excluded because at near \(r = 0\) the radial functions behaviour are \(R(r) \sim r^{-1-|\kappa|+O(\alpha^2)}\). Analogously by the same reasons \(s = -1 - \sqrt{\kappa^2 \pm a_{S,P}^2}\) must be excluded if \(a_{S,P} \ll 1\).

The solution \(s = -1 - \sqrt{\kappa^2 + a_{S,P}^2}\) must be excluded for any \(a_{S,P}\).

4. Majorana Particle Case
For Majorana particle case we must put in equation (13) and equations (19)-(26) $g_V = Q = 0$:

$$\hat{k} - g \hat{Z}_0 g_A \gamma_5 - m + a_S V_S + a_P \gamma_5 f V_P)\psi_M(k) = 0,$$

(38)

whith Majorana spinors:

$$\psi_M = \frac{1}{\sqrt{2}}(\psi + \psi^c)$$

(39)

where $\phi, \chi$ defined above in (15).

In case of right-handed heavy neutrino $Z^0 N N, W^\pm l^- \nu$ vertexes are suppressed by the smallness of light neutrino mass and therefore heavy neutrinos bound state exist predominantly due to Higgs bosons exchange ($Z^0_R$-boson interact with heavy neutrino, however if it is heavier than Higgs bosons, its contribution is negligible). Also, due to the smallness of $Z^0 N \nu, W^\pm l^- \nu$ vertexes widths of the decays $N \rightarrow Z^0_R \nu, W^\pm l^\mp$ are small in comparison with energy levels.

**5. Scalars**

Our consideration is also applicable to scalar neutrinos bound state which appear e.g. in Minimal Supersymmetric Standard Model (for Minimal Supersymmetric Standard Model see [8], [9] and references therein) $\tilde{\nu}_i \tilde{\nu}_j, \tilde{\nu}_i \tilde{\nu}_j, \tilde{\chi}_i^0 \tilde{\chi}_j^0$ bound states of scalar+ pseudoscalar.

Klein-Gordon equation of motion for scalar neutrino (antineutrino) in potentials $Z_0, V_S, P$ takes the form:

$$((E - e(g_V Z_0))^2 - \vec{k}^2 - m^2 + b_S V_S + b_P V_P)u(k) = 0,$$

(40)

At small $r$ we find the solution in the form:

$$g_i(r) = a_i r^s,$$

(41)
where only the most non-singular solution must be taken into account:

\[ s = \frac{1}{2}(-1 + \sqrt{1 - \alpha^2 (g_V g_V^\dagger)^2}) \]  

(42)

Depending on sign of \( g_V \), the equation (40) describes scalar neutrino-scalar neutrino, scalar neutrino-scalar antineutrino systems.

Now we make numerical calculation which will be described in details in our next paper [32].

7. Bethe-Salpeter equations

If \( Q_a, Q_b, T_a, T_b \) are charges and isospins of two Dirac particles \( a \) and \( b \), the Bethe-Salpeter (see e.g. [33] and references therein, for Bethe Salpeter equation for scalars systems)

equation of the such system has the following form:

\[
(\hat{k}_a - m_a)\chi(\hat{k}_b - m_b) = e^2 \int \frac{d^4k}{(2\pi)^4} \frac{(g^a_V \gamma_n + g^b_A \gamma_5)\chi(\gamma_n g^b_V + g^b_A \gamma_5)}{(p - \vec{k})^2 + M^2 + P_T} + e^2 Q_a Q_b \int \frac{d^4k}{(2\pi)^4} \frac{\gamma_n \chi}{(p - \vec{k})^2 + m_p^2 + P_T} + \sum A_s a_s a_s^n \int \frac{d^4k}{(2\pi)^4} \frac{\chi}{(p - \vec{k})^2 + m_p^2 + P_T} + \sum S_s a_s a_s^n \int \frac{d^4k}{(2\pi)^4} \frac{\chi}{(p - \vec{k})^2 + m_p^2 + P_T},
\]

(43)

where \( g^a_V = \frac{1}{c_{WW}} (T^a_V - 2Q_{sW}) \), \( g^b_A = \frac{1}{c_{WW}} 2T^a_V \).

If one of the particle is Majorana neutrino we must put in the above equation:

\[ g^a_V = Q_a = 0, \]

(44)

if \( a \) is Majorana neutrino;

\[ g^b_V = Q_b = 0, \]

(45)

if \( b \) is Majorana neutrino; and

\[ g^{a,b}_V = Q_{a,b} = 0, \]

(46)

10
if both particles are Majorana fermions.

For the same flavour particles bound states (e.g. like $l_i^-, l_i^-$) we must take into account also $u, s$-channels exchanges (see Fig.1). It must be noted that other loops with triangle vertex and four-point integrals (Fig.1) which are of order $O\left(\frac{G^2}{r^2}\right)$ at large ($r \gg \frac{1}{m_Z}$) distances, should be taken into account too.

Analogously, for scalar particles case we obtain:

\[
(k_a^2 - m_a^2)\chi(k_b^2 - m_b^2) = g_a^V g_b^V \int \frac{d^4k}{(2\pi)^4} \frac{(k_a + q)(2k_b + k_a - q)\chi}{(\vec{q} - k_a)^2 + m_Z^2 + P_T}
\]

\[
+ e^2 Q_a Q_b \int \frac{d^4k}{(2\pi)^4} \frac{\chi}{(\vec{p} - k)^2 + m_Z^2 + P_T}
\]

\[
\sum b_S^a b_S^b \int \frac{d^4k}{(2\pi)^4} \frac{\chi}{(\vec{p} - k)^2 + m_Z^2 + P_S} \tag{47}
\]

where coefficients $b_S^a$ are Yukawa couplings of scalar leptons to Higgs bosons (see e.g. [8] for Higgs bosons- scalar leptons interaction within MSSM).

In theories with R-parity violation it is possible also $\bar{\nu}l\bar{l}$ couplings are also possible. In Bethe-Salpeter equation for fermions in this case we obtain terms $\chi\gamma_5$ or $\chi\gamma_5$ which lead to $P$-parity violation which will manifested in nonrelativistic lagrangian as $V(r) \sim \sigma \vec{r}$ [8].

Acknowledgements

The authors express his sincere gratitude to all participants of the seminar which took place in Yerevan Physics Institute on October 29, 1998 and to I.Aznauryan, A.Allakhverdyan, G.Griroryan, A.Melikyan, R.Pogossyan, D.Sahakyan, G.Yegiyan, E.Prokhorenko for fruitful discussions.

References
[1] J. Feinberg, J. Sucher, Phys. Rev. 166 (1968) 1638 (see also ref. in this paper)

[2] J. Feinberg, J. Sucher, C. K. Au, Phys. Rep. 180 (1989) 83

[3] S. D. Hsu, H. P. Sikivie, Phys. Rev. 49D (1994) 4951

[4] E. Fischbach, Ann. Phys. (N.Y.) 247 (1993) 213

[5] J. A. Grifols, E. Masso, R. Toldra, Phys. Lett. B389 (1996) 186

[6] A. Yu. Smirnov, F. Vissani, hep-ph/9604443, 1996

[7] F. Ferrer, M. Nowakowski, hep-ph/9810551, 1998

[8] H. E. Haber, G. L. Kane, Phys. Rep. 117 (1985) 75

[9] J. F. Gunion, H. E. Haber, Nucl. Phys. 1986 V.B272, p.1

[10] J. C. Pati, A. Salam, Phys. Rev. D11 (1975) 566

[11] R. N. Mohapatra, J. C. Pati, Phys. Rev. D11 (1975) 566

[12] R. N. Mohapatra, G. Senjanovic, Phys. Rev. D12 (1975) 1502

[13] G. Senjanovic, Nucl. Phys. B153 (1979) 334

[14] J. F. Gunion, J. A. Grifols, A. Mendez, B. Kayser, F. Olness, Phys. Rev. D40 (1989) 1546

[15] R. N. Mohapatra, G. Senjanovic, Phys. Rev. Lett. V.44 (1980) 912

[16] R. N. Mohapatra, G. Senjanovic, Phys. Rev. D23 (1981) 165

[17] A. Kusenko, V. Kusmin, I. Tkachev, hep-ph/9801405, 1998
[18] R.A.Alanakyan, Preprint YeRPHI-1523(23)-98, 1998 "Long-Range Forces mediated by photons, gluons and fermions loops in scalars and pseudoscalars exchanges"

[19] L.B.Okun, "Leptons and Quarks"

[20] J.Steiberger, Phys.Rev.1949 V.76, P.1180

[21] J.S.Bell, R.Jackiw, Nuovo Cimento A60(1969)47

[22] P.Fayet, Phys.Lett.B69(1977)489

[23] G.Farrar, P.Fayet, Phys.Lett.B76(1978)575

[24] N.Sakai, T.Yanagida, Nucl.Phys.B197(1982)533

[25] C.S.Aulakh, R.N.Mohapatra Z.Phys.Lett.B119(1983)136

[26] F.Zwirner, Phys.Lett.B132(1983)103

[27] L.J.Hall, M.Suzuki, Nucl.Phys.B231(1984)419

[28] S.Dawson, Nucl.Phys.B261(1985)297

[29] R.Plaga, hep-ph/9610545

[30] J.A.I.Akhiezer, V.B.Beresteckij, "Quantum Electrodynamics"

[31] L.D.Landau, E.B.Lifshits, "Quantum Electrodynamics"

[32] R.A.Alanakyan, E.B.Prokhorenko, Preprint YeRPHI-1532(32)-98, 1998 (in preparation)

[33] G.C.Wick, Phys.Rev.96,(1954)1124,