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NEUTRINO MASS AND OSCILLATION ANGLE PHENOMENA WITHIN THE ASYMMETRIC LEFT-RIGHT MODELS

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ABSTRACT

The light and heavy Majorana neutrinos which appear naturally in $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model are investigated. The exact solutions are presented for the system of two neutrinos with multipole moments propagating through magnetic and matter fields. The cross section of the reaction $e^-e^- \rightarrow W_k^- W_n^-$ calculated and its dependence on the mass of the right-handed neutrino and the oscillation angle is investigated. The process $e^+e^- \rightarrow W_k^+ W_n^-$ is also included in our analysis.

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1 Introduction

A precise knowledge of the intrinsic properties of the neutrino would have a large impact on our understanding of both astrophysics and particle physics. At present the experimental values of neutrino masses, lifetimes, decay modes, the oscillation angles and multipole moments (MM) are still inconclusive. In addition the question remains open as to whether the neutrino is a Majorana or Dirac particle. The standard model (SM) predicts three massless neutrinos while in almost every extension of the SM there are more than three massive neutrinos.

In this paper we will draw attention to the possible manifestations of the left-right (LR) symmetry from the point of view of the experiments with cosmic neutrinos and the collider experiments going through neutrino-exchange. We shall compare the results of the standard model (SM) with those obtained within the asymmetric LR model (ALRM). Our results will be represented for the ALRM proposed in Ref.[1]. This model has one more additional parameter φ named by the LR symmetry violating angle. Varying φ one may obtain all the possible versions of the LR models. The SM is reproduced in the SM particles sector at $\varphi = 0$ and when the following conditions are fulfilled

$$\Phi = \xi = 0 \quad (1)$$

where Φ and ξ are the mixing angles of the neutral and charged gauge bosons, respectively, and

$$g_L = e s_W, \quad g'^{-1} = \sqrt{c_W^2 e^{-2} - g_R^{-2}}, \quad |g_R| > g_L s_W c_W^{-1} \quad (2)$$

where g_L , g_R and g' are the couplings constants (CC's) of $SU(2)_L$, $SU(2)_R$ and $U(1)_{B-L}$ gauge groups, respectively, and $c_W = \cos \theta_W$, $s_W = \sin \theta_W$. In any LR model there are light left-handed and heavy right-handed neutrinos in each lepton generation. Depending on the choice of the Higgs sector the neutrinos could be either Dirac or Majorana particles. The heavy and the light neutrinos of the same flavour are obtained from mixing the mass eigenstates ν_i and N_i according to

$$\nu_{iL} = \nu_i \cos \theta_{\nu_i N_i} + N_i \sin \theta_{\nu_i N_i} \quad (3)$$

$$\nu_{iR} = -\nu_i \sin \theta_{\nu_i N_i} + N_i \cos \theta_{\nu_i N_i} \quad (4)$$

where $\theta_{\nu_i N_i}$ is the oscillation angle in the vacuum. Mixing between different generations is also allowed. The present limit on the mass of the electron neutrino is [2]

$$m_{\nu_e} < 8 \text{ eV} \quad (5)$$

whereas for its heavy counterpart N_i the following mass bounds have been derived [3,4]

$$(377 - 413) \text{ GeV} \left(\frac{1 \text{ TeV}}{m_{W_2}} \right)^4 < m_{N_e} < 2e^{-1} \sqrt{\sin \theta_W (m_{W_1}^2 + m_{W_2}^2)} \quad (6)$$

where m_{W_i} ($i = 1, 2$) denote the masses of the left- and right-handed W-bosons, respectively, and the numbers in parentheses reflects the uncertainty from the nuclear matrix element of the neutrinoless double β -decay

$$^{76}\text{Ge} \rightarrow ^{76}\text{Se} e^+ e^- \quad (7)$$

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on which the lower bound estimate is based. Information about the oscillation parameters could be obtained from accelerator and nuclear reactor experiments and from available data on atmospheric and solar neutrinos (for Review see [5]).

The aim of this paper is to consider the influence of neutrino parameters on neutrino oscillations and $e^-e^-(e^+)$ collider- and very high energy cosmic ray experiments. The paper is organized as follows: In Section II we investigate the Majorana neutrinos propagating through magnetic and dense matter fields. Section III is devoted to the investigation of the indirect signatures of the neutrino parameters in the total cross section of the reactions

$$e^-e^- \longrightarrow W_k^- W_n^- \quad (k, n = 1, 2) \quad (8)$$

$$e^+e^- \longrightarrow W_k^+ W_n^- \quad (9)$$

Section IV summarizes our conclusions.

2 Neutrino oscillations

In this section we consider the propagation of massive Majorana neutrinos through matter and a magnetic field taking into account their electromagnetic properties. We have in mind neutrinos emanating from the inner core of a star which on their way out have to traverse a sizeable magnetic field. The electromagnetic quantities we are interested in are caused by radiative corrections which at the one loop level have their origin in two kinds of diagrams shown in Fig.1, where $V = \gamma, Z_{1,2}$. The subject has already been addressed by Voloshin, Vysotsky and Okun in Ref.[6], however, only in the framework of the symmetric left-right model assuming Dirac nature of neutrinos and disregarding right-handed charged current neutrino-electron interaction. Within the SM Majorana neutrinos have been discussed in Ref.[7,8].

We rely our investigation on the most general form of the left-right symmetric gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [1] and taking into account right-handed charged and neutral currents. The basic interaction including the processes defined in Figs.1 a,b can be formulated in terms of the following effective Lagrangians

$$\mathcal{L}_a = \frac{1}{2} a_{L,R} \bar{\nu}_{L,R} \partial_\mu V_{\mu\nu} \gamma_5 \nu_{L,R} + \text{conj.} \quad (10)$$

$$\mathcal{L}_b = \frac{1}{2} \mu_{L,R} \bar{\nu}_{L,R} \sigma_{\mu\nu} V_{\mu\nu} \nu_{L,R} + \text{conj} \quad (11)$$

where $V_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu$. These give rise to the appearance of so called anapole electromagnetic (weak) moments [9] described by the coefficients $a_{L,R}$ (which might be different for every lepton family, e. g. $a_{L,R}^\mu \neq a_{L,R}^\nu$) and electromagnetic (weak) dipole moments described by $\mu_{L,R}$. We should stress that in the one-flavour approximation for the Majorana neutrino the anapole moment is the only electromagnetic form factor allowed by the CTP -invariance. For massless neutrinos it is connected with the neutrino electric charge radius (ECR). In the framework of the SM the ECR has been calculated in Ref. [10]. The results display the dependence on the top quark and the Higgs boson masses and the neutrino flavour. For example, at $m_t = 90 \text{ GeV}$ and $m_H = 100 \text{ GeV}$ one obtains

$$\langle r^2 \rangle \cong -(13.4 \pm 3.6) \cdot (10^{-17} \text{ cm})^2 \quad (12)$$

Within the SM the influence of the ECR and anapole moments of neutrinos has been considered for the disintegration of deuteron by neutrinos [11]. It has been shown that the ECR contribution can be as large as $\approx 1.2\%$. Hence, these quantities should be taken into account when investigating the neutrino phenomena.

The appearance of a neutrino dipole moment is also possible in the SM but its value is very small

$$(\mu_\nu)_{SM} \approx 2.7 \cdot 10^{-10} \mu_B \frac{m_{\nu_l}}{m_p} \quad (13)$$

where m_{ν_l} and m_p are the masses of the neutrino and the proton, respectively, and μ_B is Bohr magneton. Such a dipole moment could be much larger in ALRM where it is proportional

$$\mu_\nu \approx \frac{3.6 \cdot 10^{-10} g_R m_l \mu_B \sin 2\xi}{g_L m_p} \quad (14)$$

The current limits on magnetic dipole moment are coming from reactor experiments and cosmological observations. There is one bound on the electron neutrino moment inferred from laboratory experiments of neutrino-electron scattering and $e^+e^- \longrightarrow \gamma \nu \nu$ process [12] giving $\mu_{\nu_e} < 10.8 \cdot 10^{-10} \mu_B$. Astrophysical considerations have led to a somewhat more stringent upper bound [13] $\mu_{\nu_e} < 7.9 \cdot 10^{-10} \mu_B$. In the literature the possibility of measuring μ_{ν_e} by use of the transition radiation of neutrinos crossing the interface between two media [14] is also discussed. For the Dirac particle the values of the anapole moment could be estimated by comparison with the diagonal elements of the magnetic moment $\mu_\parallel$. In the case of Majorana neutrinos $\mu_\parallel$ are equal to zero and only the off diagonal elements (transition magnetic moments) can exist.

First we consider the case when a light left-handed electron neutrino ν_L can mix with a heavy right-handed electron neutrino ν_R . Then their connections with the mass eigenstates ν and N are given by (3) and (4). Next we shall take into account the coherent interaction of the neutrino with matter and assume the matter to be electric neutral. Then by the standard way from the Lagrangian averaged over the matter distribution we obtain the neutrino transmission equation in a Schrödinger-like form

$$i \frac{d}{dz} \begin{pmatrix} \nu_R \\ \nu_L \end{pmatrix} = \mathcal{H} \begin{pmatrix} \nu_R \\ \nu_L \end{pmatrix} \quad (15)$$

where

$$\begin{aligned} \mathcal{H} &= \Delta \mathcal{M} + C_L P_- + C_R P_+ - a_L P_+ (\vec{\sigma} \cdot \text{rot} \vec{H}) P_+ \\ &\quad + a_R P_- (\vec{\sigma} \cdot \text{rot} \vec{H}) P_- - \mu P_+ (\vec{\sigma} \cdot \vec{H}) P_- - \mu P_- (\vec{\sigma} \cdot \vec{H}) P_+ \\ C_L &= \sqrt{2} G_F n_e + \frac{n_n}{4} \sum_{i=1}^2 \frac{g_{V_i}^L}{m_{Z_i}^2} (g_{V_i}^d + g_{A_i}^d) \\ C_R &= \frac{g_R^2 n_e}{4} \left(\frac{\sin^2 \xi}{m_{W_1}^2} + \frac{\cos^2 \xi}{m_{W_2}^2} \right) + \frac{n_n}{4} \sum_{i=1}^2 \frac{g_{V_i}^R}{m_{Z_i}^2} (g_{V_i}^d - g_{A_i}^d) \\ \Delta \mathcal{M} &= \begin{pmatrix} \frac{\Delta m^2}{4E} \cos 2\theta_{\nu_e N_e} & -\frac{\Delta m^2}{4E} \sin 2\theta_{\nu_e N_e} \\ -\frac{\Delta m^2}{4E} \sin 2\theta_{\nu_e N_e} & -\frac{\Delta m^2}{4E} \cos 2\theta_{\nu_e N_e} \end{pmatrix} \end{aligned}$$

n_e and n_p are the densities of electrons and neutrons, respectively. The operator $P_\pm = \frac{1}{2}(1 \pm \sigma_3)$ projects on ν_L and ν_R , respectively. We choose the propagation direction of the

neutrino as the axis of the spin quantization (hereafter we use the notations of Ref.[1]). We neglect irrelevant terms in the Hamiltonian which lead only to an unobservable phase. The solution of (15) could be found in the realistic case when we use the values of a density and magnetic field profile for the star interior. Usually, in order to illustrate the general properties of the neutrino system the equation is solved in the case of matter with a constant density along z -axis, $H_\perp \neq H_\perp(z)$, where $H_x + iH_y = H_\perp \exp i\phi(z)$, and steady rotation of $H_\perp$ (the geometrical phase velocity $\frac{d\phi(z)}{dz} \equiv \dot{\phi} = \text{const}$). Then the solution can be written as

$$\nu(z) = e^{(-i\vec{\tau}\vec{\sigma}z)}\nu(0) = (\cos\omega z - i(\vec{\tau}\vec{\sigma})\omega^{-1}\sin\omega z)\nu(0) \quad (16)$$

where

$$\nu(z) = \begin{pmatrix} \nu_R \\ \nu_L \end{pmatrix}, \quad \vec{\tau} = (-\mu H_x - \Delta\mathcal{M}_{12}, \mu H_y, \frac{C_R - C_L}{2} + \frac{a_R + a_L}{2}(\text{rot}\vec{H})_z + \Delta\mathcal{M}_{11} + \frac{\dot{\phi}}{2})$$

and $\omega = \sqrt{\vec{\tau}^2}$. When the initial state is purely left-handed, then the numbers of the left-handed and the right-handed neutrinos at the point z are defined by the expressions

$$\mathcal{N}_L = 1 - \frac{1}{2}\sin^2 2\theta_m (1 - \cos \frac{2\pi\omega z}{\pi}) \quad (17)$$

$$\mathcal{N}_R = \frac{1}{2}\sin^2 2\theta_m (1 - \cos \frac{2\pi\omega z}{\pi}) \quad (18)$$

where θ_m is the oscillation angle in the matter and the magnetic field

$$\sin^2 2\theta_m = \frac{\vec{\tau}_\perp^2}{\omega^2} \quad (19)$$

As follows from (17) and (18) $\frac{\pi}{\omega}$ is the oscillation length. The results of the SM follow when we set

$$\varphi = \Phi = \xi = g_R = 0, \quad m_{Z_2} = \infty \quad (20)$$

The Mikheyev-Smirnov-Wolfenstein (MSW) effect takes place at the condition

$$C_L(\epsilon - 1) + (a_R + a_L)(\text{rot}\vec{H})_z + 2\Delta\mathcal{M}_{11} + \dot{\phi} = 0 \quad (21)$$

where $\epsilon = \frac{C_R}{C_L}$. It defines the resonance density at which the mixing angle θ_m becomes $\frac{\pi}{4}$. We see that the appearance of the anapole moment modifies the resonance condition. It is obvious that using the solar neutrino data we could probe the magnetic field structure of the sun generated by the so-called dynamo mechanism. According to this mechanism the strength of the magnetic field at the bottom of the convection zone is constrained to a value between 40 and 150 kG [15]. The sun's magnetic field is non-potential, namely

$$(\text{rot}\vec{H})_z = J_z \neq 0 \quad (22)$$

The levels of J_z is about $1.5 \cdot 10^{-3} \text{ Am}^{-2}$. The longitudinal electric currents could flow in opposite directions and change their directions after a solar flare [16]. We should also like

to remark that the condition (21) can be achieved in different ways. For example, the second and the fourth terms in (21) have to cancel if $C_L(1 - \epsilon) = 2\Delta\mathcal{M}_{11}$. From (19) it follows when the effects of matter and geometrical phase are neglected, i.e. $C_L = C_R = \dot{\phi} = 0$, the effects of the mass difference and the magnetic field have the same order of magnitude in the case $\Delta m^2 \leq 10^{-7} \text{ eV}^2$ at $H_\perp \sim 10^3 \text{ G}$, $E \sim 10 \text{ MeV}$ and $\mu = 10^{-10} \mu_B$. However, when the matter, the mass difference and the geometrical phase velocity terms are cancelled among themselves the effects of spin flip caused by the magnetic field would be dominant and could be correlated with sunspot activity. Obviously for $(m_{N_e}^2 - m_{\nu_e}^2) > 10^{-4} \text{ eV}^2$, the solar neutrino puzzle suggests contribution of $\nu_e - \nu_\chi$ ($\chi = \mu, \tau$) oscillations, too. We recall that the current bounds on light left-handed neutrinos are [17]

$$\begin{aligned} m_{\nu_\mu} &< 270 \text{ keV} \\ m_{\nu_\tau} &< 31 \text{ MeV} \end{aligned} \quad (23)$$

The expressions for the probabilities to find ν_{eL} and $\nu_{\chi L}$ follow from (17) and (18) by the substitution

$$C_R \rightarrow \frac{n_n}{4} \sum_{i=1}^2 \frac{g_{V_i}^{\nu_L}}{m_{Z_i}^2} (g_{V_i}^d + g_{A_i}^d), \quad \mu_{\nu_e N_e} \rightarrow \mu_{\nu_e \nu_\chi}, \quad a_{Re} = a_{L\chi} \quad (24)$$

It should be noted here that if $\nu_e - \nu_\chi$ oscillations exist the following reaction must be observed $e^- \chi^+ \rightarrow W^+ W^-$. This reaction could be observable with high-energy cosmic rays ($E_{\text{max}} \sim 10^{20} \text{ eV}$).

Now we can investigate the energy-level scheme of the four, say ν_{eR} , ν_{eL} , $\nu_{\chi R}$ and $\nu_{\chi L}$ neutrino chiral states for the sun, the supernova and the neutron stars. Equating the diagonal elements of the Hamiltonian which define the energies of the flavour levels of the neutrino system we obtain the energy-level crossings for six channels. Our analysis shows the resonance picture is similar to that predicted by the SM [8]. Of course, depending on the choice of the structural parameters (SP's) of the ALRM the location and the distances between resonances are changed. We could also face the case of resonances permutations. Here the main question we are interested in is how big the deviations from the SM could be. For example, choosing $g_R = 1.6g_L$, for the channel $\nu_{eR} - \nu_{eL}$ in the convection zone and the upper radiation zone where $n_e \cong 6n_n$ we obtain $\epsilon \cong 10\%$. In our estimation we have used the following values of the SP's of the ALRM

$$\begin{aligned} m_{W_1} &= 80.13 \text{ GeV}, \quad m_{W_2} = 477 \text{ GeV}, \quad m_{Z_1} = 91.77 \text{ GeV}, \quad m_{Z_2} = 800 \text{ GeV}, \\ \Phi &= 9.6 \cdot 10^{-3}, \quad \xi = 3.1 \cdot 10^{-2}, \quad \theta_{\nu N_i} = 5 \cdot 10^{-2} \end{aligned} \quad (25)$$

which are in accordance with the existing experimental bounds (see e.g. Ref. [18]). Hereafter we assume that $\varphi = 0$ and the CC's are defined by (2) and $g' > 0$.

In conclusion in this case we note that in order to carry out the realistic calculation one needs to know the anapole moment values of the neutrino. It should be also kept in mind that the neutrinos are perfect for testing the equivalence principle [19] since they are subject only to weak and gravitational forces. However, even if this principle is not violated one should include the effects of the gravitational field (see e. g. Ref. [20]). We are going to investigate these problems in our forthcoming work.

3 Experiments with electron and positron beams

In this section we discuss the impact of the Majorana neutrino scenario both for collider experiments and for experiments with cosmic rays. In our calculations we shall use the Feynman rules for the case of Majorana neutrinos which has been formulated in Ref. [21]. As it follows from them the cross sections of lepton flavor conserving processes have the same form both for the Dirac and Majorana neutrinos.

A reliable test of the nature of neutrino involved is the process which is connected (by time reversal symmetry) to neutrinoless double β decay, i.e.

$$e^- e^- \longrightarrow W_k^- W_n^- \quad (k, n = 1, 2) \quad (26)$$

We remember that it is forbidden if the neutrino is a Dirac particle. The reaction (26), represented by the Feynman diagrams of Fig.2, will be accessible to cosmic rays and to the $e^- e^-$ option of LEP II and NLC (TLC). The experimental program is motivated not only by looking for lepton flavor violation but also by searching for doubly charged leptons and Higgs particles.

The modification of the Higgs sector in the SM could also leads to the appearance of the heavy Majorana neutrino which is mixed with the light one. Then, in the second order of the perturbation theory the process (26) proceeds via neutrino-exchanges in t - and u -channels. As only weak CC's take part in the process we could expect the violation of the unitarity bound for the cross section. Really in the high-energy limit $s \rightarrow \infty$ where s is the energy in the center of mass system, the SM total cross section has the form [22]

$$\sigma \sim \frac{G_F^2 m_N^2}{4\pi} \sin^4 \alpha \quad (27)$$

where $\alpha = \theta_{\nu_e N_e}$. The analogical situation could also occur for the processes caused by the electromagnetic CC's only. For example, the total cross section of the Compton effect on W -boson in the ultrarelativistic case is defined by

$$\sigma \sim \frac{5\pi e^4 E_\gamma}{576\pi^2 m_W^2} \quad (28)$$

where E_γ is the energy of the incident photon. Of course, in the high orders of the perturbation theory, when all the CC's of the electroweak theory come into play, the unitarity bound is not a problem.

Within the symmetric LRM the process (26) has been studied in Refs.[22-24]. In Refs.[22,24] it was assumed that the reaction (26) was caused by the neutrino-exchange only. In Ref.[23] the model was used in which the process (26) proceeds not only via a neutrino-exchange in t -channel but via a doubly charged triplet Higgs Δ^{--} -exchange in s -channel. It should be noted that the presence of the diagram with Δ^{--} leads to the cancellation in the cross section of the terms which violate the unitarity.

The conclusions of Refs.[22] and [24] do not coincide with each other. In Ref.[22] in the high-energy limit the total cross section grows as a linear function of s , while in Ref.[24] it increases as the square of s .

Let us calculate the cross section of the process (26) within the ALRM. The main difference of this model from the SM is that if in the latter the reaction (26) occurs for left-polarized electrons ($e_L^- e_L^-$) only, then in the ALRM the contribution to the cross

section is given by both $e_{L,R}^- e_{L,R}^-$ and $e_{L,R}^- e_{R,L}^-$ beams. It is worth reminding that the expression for the process

$$e_L^- e_L^- \longrightarrow W_1^- W_1^- \quad (29)$$

obtained in the ALRM should coincide with that in the SM.

The following identity

$$\bar{e} \gamma_\mu (1 \pm \gamma_5) \nu_{L,R} = -\bar{\nu}_{L,R} \gamma_\mu (1 \mp \gamma_5) e^c \quad (30)$$

will be useful in our calculations. For example, with its help the matrix element corresponding to the process

$$e_L^-(k_1) e_R^-(k_2) \longrightarrow W_k^-(p_1) W_n^-(p_2) \quad (31)$$

takes the form

$$M_{LR} = \frac{1}{2} (\sin 2\alpha) g_L g_R b_L^{(k)} b_R^{(n)} \bar{e}(k_1) \gamma_\mu (1 + \gamma_5) \left[\frac{(\hat{k}_1 - \hat{p}_1) + m_\nu}{(k_1 - p_1)^2 - m_\nu^2} - \frac{(\hat{k}_1 - \hat{p}_1) + m_N}{(k_1 - p_1)^2 - m_N^2} \right] \gamma_\lambda (1 + \gamma_5) e^c(k_2) W_\mu^{*(k)}(p_1) W_\lambda^{*(n)}(p_2) + (k \leftrightarrow n) \quad (32)$$

where $b_{L,R}^{(k)} = \frac{1}{2} \{ [\pm 1 + (-1)^k] \sin \xi + [1 \mp (-1)^k] \cos \xi \}$. With a precision up to constant this expression coincides with the amplitude of t -channel of the reaction

$$f_i \bar{f}_i \longrightarrow W_k^- Z_n \quad (k, n = 1, 2) \quad (33)$$

where i is the flavour of a fermion f , which within the SM ($k = n = 1$) was investigated in Ref. [25] whereas within the LRM it was done in Ref. [26]. Therefore we can use the results of these works and write down the differential cross section in the following form

$$\frac{d\sigma_{L,R}^{(kn)}}{dt} = \frac{(g_L g_R m_N^2 \sin 2\alpha)^2}{512\pi s^2} \left[\left(\frac{b_R^{(n)} b_L^{(k)}}{t - m_N^2} \right)^2 B_3^{(kn)}(s, t, u) + \left(\frac{b_R^{(k)} b_L^{(n)}}{u - m_N^2} \right)^2 B_3^{(kn)}(s, u, t) + \frac{b_L^{(k)} b_L^{(n)} b_R^{(k)} b_R^{(n)}}{(t - m_N^2)(u - m_N^2)} B_4^{(kn)}(s, t, u) \right] \quad (34)$$

where the functions $B_{3,4}^{(kn)}(s, u, t)$ are defined in Ref.[26]. Hereafter we for the sake of simplicity take the mass of the light neutrino to be equal zero.

The analogical calculations for the differential cross section of the reaction (29) give the expression

$$\frac{d\sigma_L^{(kn)}}{dt} = \frac{(g_L^2 m_N b_L^{(k)} b_L^{(n)} \sin^2 \alpha)^2}{512\pi s^2} \left\{ \left[3s - \frac{Fs^2(s - 2m_{W_n}^2 - 2m_{W_k}^2)}{4m_{W_n}^2 m_{W_k}^2} - \frac{st(s - m_{W_n}^2 - m_{W_k}^2)^2}{m_{W_n}^2 m_{W_k}^2} \right] \frac{1}{(t - m_N^2)^2} + (u \leftrightarrow t) + [2s + \frac{Fs^2(s - 2m_{W_n}^2 - 2m_{W_k}^2)}{2m_{W_n}^2 m_{W_k}^2}] \frac{1}{(t - m_N^2)(u - m_N^2)} \right\} \quad (35)$$

where $F = 4(ut - m_{W_n}^2 m_{W_k}^2)/s^2$. The differential cross section of the reaction

$$e_R^- e_R^- \longrightarrow W_k^- W_n^- \quad (36)$$

follows from (35) by the substitution

$$g_L \rightarrow g_R, \quad b_L^{(k)} \rightarrow b_R^{(k)}, \quad \alpha \rightarrow \alpha + \frac{\pi}{2} \quad (37)$$

Now, using the obtained expressions we can evaluate the differential cross section of the reaction (26) for the case of the polarized electrons and unpolarized W -bosons. Because of the t and u -poles it displays a strong angular dependence. This circumstance allows to distinguish the model under investigation from the one, where the Δ^{--} -exchange in s -channel is dominant. In the latter case we shall have an isotropic angular distribution.

The total cross section of the reaction (31) follows from (34) after multiplication by $\pi\beta s$ and substitution

$$\left. \begin{aligned} \frac{B_3^{(kn)}(s, t, u)}{(t - m_N^2)^2} \\ \frac{B_3^{(kn)}(s, u, t)}{(u - m_N^2)^2} \end{aligned} \right\} \rightarrow D_3^{(kn)}(s), \quad \frac{B_4^{(kn)}(s, t, u)}{(t - m_N^2)(u - m_N^2)} \rightarrow D_4^{(kn)}(s)$$

where

$$\begin{aligned} D_3^{(kn)}(s) &= \frac{1}{m_{W_n}^2 m_{W_k}^2} \left[1 + \frac{(m_{W_n}^2 + m_{W_k}^2 - s - 2m_N^2) \ln L_\nu}{2\beta s} \right] \\ &+ \frac{1}{m_N^4} \left\{ \frac{[4m_{W_n}^2 m_{W_k}^2 - 2m_N^2(m_{W_n}^2 + m_{W_k}^2 - s)](\ln L_\nu - \ln L)}{\beta s m_N^2} \right. \\ &\left. - 4 \right\} - \frac{s(m_{W_n}^2 + m_{W_k}^2)}{C_\nu m_{W_n}^2 m_{W_k}^2}, \\ D_4^{(kn)}(s) &= - \frac{2[2s(m_{W_n}^2 + m_{W_k}^2) + C_\nu] \ln L_\nu}{\beta s m_{W_n}^2 m_{W_k}^2 (m_{W_n}^2 + m_{W_k}^2 - s - 2m_N^2)} \\ &- \frac{1}{m_{W_n}^2 m_{W_k}^2} - \frac{8(m_{W_n}^2 + m_{W_k}^2)}{\beta m_N^2 (m_N^2 - m_{W_n}^2 - m_{W_k}^2 + s)} \left(\frac{\ln L_\nu}{m_{W_n}^2 + m_{W_k}^2 - s - 2m_N^2} \right. \\ &\left. - \frac{\ln L}{m_{W_n}^2 + m_{W_k}^2 - s} \right), \\ C_\nu &= m_N^2 (m_{W_n}^2 + m_{W_k}^2 - s - m_N^2) - m_{W_n}^2 m_{W_k}^2 \\ L_\nu &= \left| \frac{m_{W_n}^2 + m_{W_k}^2 - s - 2m_N^2 + \beta s}{m_{W_n}^2 + m_{W_k}^2 - s - 2m_N^2 - \beta s} \right|, \\ L &= L_\nu |_{m_N=0}, \quad \beta = \sqrt{\left(1 - \frac{m_{W_n}^2 + m_{W_k}^2}{s}\right)^2 - \left(\frac{2m_{W_n} m_{W_k}}{s}\right)^2} \end{aligned}$$

For the total cross section of the reaction (29) we find the expression

$$\sigma_L^{(kn)} = \beta \frac{(g_L^2 m_N b_L^{(k)} b_L^{(n)} \sin^2 \alpha)^2}{256\pi s} \left\{ \frac{2s - 4(m_{W_n}^2 + m_{W_k}^2)}{m_{W_n}^2 m_{W_k}^2} - s \left(3 + \frac{m_N^4 + m_{W_n}^2 m_{W_k}^2}{m_{W_n}^2 m_{W_k}^2} \right) C_\nu^{-1} \right.$$

$$\begin{aligned} &+ \frac{\ln L_\nu}{\beta s} \left\{ \frac{2sm_N^2 + 2(m_{W_n}^2 + m_{W_k}^2 - s - 2m_N^2)(m_{W_n}^2 + m_{W_k}^2)}{m_{W_n}^2 m_{W_k}^2} \right. \\ &\left. + \frac{2}{m_{W_n}^2 + m_{W_k}^2 - s - 2m_N^2} \left(s + \frac{s - 2m_{W_n}^2 - 2m_{W_k}^2}{m_{W_n}^2 m_{W_k}^2} C_\nu \right) \right\} \end{aligned} \quad (38)$$

Having done the replacement (37) in the expression (38) we obtain the total cross section of the reaction (36). In the high energy limit from (38) it follows the formulae (27). In contrast, $\sigma_{LR}^{(kn)}$ has a good asymptotic behavior

$$\sigma_{LR}^{(kn)} \sim \frac{\ln s}{s} \quad (39)$$

It is due to the fact, that in the reactions (29) and (36) the virtual neutrino flips the helicity between the acts of its emission and its absorption. As a result their amplitudes are proportional to the neutrino masses. As flipping the neutrino helicity does not take place in the reaction (31) its amplitude contains the two terms with the opposite signs. They are caused by ν - and N - exchanges, respectively. As the calculations show, the disregard of the contribution of the light neutrino would lead to the cross section being proportional to s . We remind that in Ref.[23] the cross section of the reaction (31) was calculated at the neglect of the ν -exchange.

As it follows from the expression for the total cross section of the reaction (26) $\sigma^{(kn)}$ this reaction could be a good tool for the definition of such parameters of the ALRM as m_N , g_R , α and ξ . In our numerical calculation we shall, in what follows, use the values of the SP's of the ALRM defined by (25). The values of σ^{kn} strongly depends on m_N . For example, at $m_N = 100 \text{ GeV}$ and $g_L = g_R$, it is about 10^{-2} fb . In this case $\sigma_{LR}^{(11)}$ and $\sigma_R^{(11)}$ have the same order of magnitude $\sim a \text{ few} \cdot 10^{-3} \text{ fb}$, whereas $\sigma_L^{(11)}$ is about $a \text{ few} \cdot 10^{-2} \text{ fb}$. At increasing m_N the contribution from $\sigma_{LR}^{(11)}$ becomes dominant. For example, at $m_N = 1 \text{ TeV}$ and $\sqrt{s} = 200 \text{ GeV}$ we have

$$\begin{aligned} \sigma_R^{(11)} &\approx 0.23 \cdot 10^{-2} \sigma_{LR}^{(11)} \\ \sigma_L^{(11)} &\approx 1.8 \cdot 10^{-2} \sigma_{LR}^{(11)} \end{aligned} \quad (40)$$

It should be noted here that at the sufficiently great m_N ($m_N > 5 \text{ TeV}$) the cross section of the reaction (26) could reach the values which are compatible and even greater than the cross section of the reaction

$$e^- e^+ \longrightarrow W_1^- W_1^+ \quad (41)$$

In Fig.3 we represent $\sigma_{LR}^{(11)}$ versus $\sqrt{s}$ for $m_N = 1$ and 1.4 TeV ($g_L = g_R$).

When $k = 1$ and $n = 2$ the main contributions into the total cross section are given by $\sigma_R^{(12)}$ and $\sigma_{LR}^{(12)}$. At $m_N = 100 \text{ GeV}$ these quantities have the same order of magnitude ($\sim a \text{ few} \cdot 10^{-2} \text{ fb}$). Again with the increase of m_N the contribution from $\sigma_{LR}^{(12)}$ gains dominance. For example, at $m_N = 1 \text{ TeV}$, $\sigma_{LR}^{(12)}$ exceeds $\sigma_R^{(12)}$ on the order of magnitude. In Fig.4 we give $\sigma_{LR}^{(12)}$ as a function of $\sqrt{s}$ for $m_N = 1$ and 1.4 TeV ($g_L = g_R$).

In the case $k = n = 2$ the reaction (26) mainly goes for the $e_R^- e_R^-$ beams. Only at the enormous values of m_N ($\sim 100 \text{ TeV}$), $\sigma_{LR}^{(22)}$ may be compatible with $\sigma_R^{(22)}$. In Fig.5 we plot $\sigma_R^{(22)}$ as a function of $\sqrt{s}$ for $m_N = 1$ and 1.4 TeV ($g_L = g_R$).

It is also important to recall, that the dependence on m_N for $\sigma_R^{(22)}$ is much weaker than the one for $\sigma_{LR}^{(22)}$. For example, varying m_N from 10^3 GeV up to 10^2 GeV leads to the decrease of $\sigma_R^{(22)}$ and $\sigma_{LR}^{(22)}$ on one and three orders of magnitude, respectively.

Now we should pay some attention to the question of detectability of the processes

$$e^-e^- \rightarrow W_k^- W_n^- \rightarrow l_i^- l_j^- + \text{neutrinos} \quad (42)$$

$$\rightarrow l_i^- + \text{jet} + \text{neutrinos} \quad (43)$$

where i and j are the fermion flavours. The most clean signature is the case $i, j = e, \mu$ plus missing momentum carried away by neutrinos. In order to eliminate the major background coming from RC to elastic e^-e^- scattering we should demand $i \neq j$ or the sizeable $p_\perp$ of outgoing leptons. The cut on $p_\perp$ helps to reduce the QED background caused by radiative pair production

$$e^-e^- \rightarrow e^-e^-\mu^+\mu^+\tau^+\tau^+ \quad (44)$$

Another relevant background processes are

$$e^-e^- \rightarrow e^-e^-Z_k \rightarrow e^-e^-l_i^-l_i^+ \quad (45)$$

$$\rightarrow e^-e^-\nu_l\bar{\nu}_l \quad (46)$$

and

$$e^-e^- \rightarrow e^-e^-W_k^-W_n^+ \rightarrow e^-e^-l_i^-l_j^+\nu_i\bar{\nu}_j \quad (47)$$

The main difference between the final states of the reactions(45)-(47) and (42)-(43) is that in the former case they consist of like-sign leptons pair only. Therefore, it is easy to distinguish the W^-W^- signal from these background by means of the charged final-state leptons identification. We also have the background caused by multihadronic events

$$e^-e^- \rightarrow e^-e^-W_k^-W_n^+q_i\bar{q}_i \quad (48)$$

For this reaction there is no possibility to observe the high $p_\perp$ of the final like-sign non-electron leptons. Hence, we conclude that the cut on $p_\perp$ and the charged lepton identification of the outgoing leptons will be very useful in order to reduce backgrounds for the $W_k^-W_n^-$ signal.

Now let us consider the process

$$e^+e^- \rightarrow W_k^-W_n^+ \quad (k, n = 1, 2) \quad (49)$$

The cross section of the reaction (49) does not depend on the assumed (Dirac or Majorana) nature of the neutrino. Unlike the reaction (26) this reaction proceeds through the s - and t -channels. All the couplings of the electroweak interaction will enter here. As a result the violating unitarity terms which enter the partial cross sections are cancelled themselves and the asymptotic behaviour of the total cross section is defined by the expression (39). The advantage of the reaction (49) compared with (26) is the possibility to measure the trilinear gauge boson couplings (electromagnetic and weak). The investigation of the weak trilinear gauge boson couplings can be studied for $k \neq n$ because in this case the channel connected with the γ -exchange is closed. At $\alpha = 0$

the reaction (49) has been considered in [1b] where one was shown that the total cross section $\sigma_{\lambda\bar{\lambda}}^{(kn)}$ where λ and $\bar{\lambda}$ are the helicities of the electron and positron, respectively, is sensitive to the m_N variations. So we shall also concentrate on $\sigma_{\lambda\bar{\lambda}}^{(kn)}$. The expression for the total cross section of the reaction (49) follows from the corresponding one of Ref.[1b] by the following substitutions

$$\begin{aligned} \mathcal{D}_2^{(kn)} &\rightarrow \mathcal{D}_2^{(kn)} \cos^2 \alpha + \mathcal{D}_{2\nu}^{(kn)} \sin^2 \alpha \\ \mathcal{D}_{2\nu}^{(kn)} &\rightarrow \mathcal{D}_{2\nu}^{(kn)} \cos^2 \alpha + \mathcal{D}_2^{(kn)} \sin^2 \alpha \\ \mathcal{D}_{3\nu}^{(kn)} &\rightarrow \mathcal{D}_{3R}^{(kn)}(\alpha) \\ \mathcal{D}_3^{(kn)} &\rightarrow \mathcal{D}_{3L}^{(kn)}(\alpha) \end{aligned} \quad (50)$$

where

$$\begin{aligned} \mathcal{D}_{3R}^{(kn)}(\alpha) &= \frac{s}{m_{W_n}^2 m_{W_k}^2} (m_{W_n}^2 + m_{W_k}^2 + \frac{\beta^2 s}{12}) - 4 + \frac{2 \ln L}{\beta s} \{ (m_{W_n}^2 + m_{W_k}^2 - s) [1 - \cos^2 \alpha (2 - \cos^2 \alpha)] + \frac{m_{W_n}^2 m_{W_k}^2 \sin^2 2\alpha}{2m_N^2} \} \\ &+ \sin^2 2\alpha + \frac{m_N^2 \cos^2 \alpha}{m_{W_n}^2 m_{W_k}^2} \left[\frac{m_{W_n}^2 + m_{W_k}^2 - s}{2} - \frac{sm_N^2 (m_{W_n}^2 + m_{W_k}^2) \cos^2 \alpha}{C_\nu} \right] \\ &- m_N^2 (1 + \cos^2 \alpha) + \frac{m_N^2 \cos^2 \alpha \ln L_\nu}{m_{W_n}^2 m_{W_k}^2 \beta s} [C_\nu (1 + \frac{4m_{W_n}^2 m_{W_k}^2 \sin^2 \alpha}{m_N^4})] \\ &+ \frac{1}{2} \cos^2 \alpha (m_N^2 + \frac{4m_{W_n}^2 m_{W_k}^2}{m_N^2}) (m_{W_n}^2 + m_{W_k}^2 - s - 2m_N^2) \\ &+ 2s(m_{W_n}^2 + m_{W_k}^2), \\ \mathcal{D}_{3L}^{(kn)}(\alpha) &= \mathcal{D}_{3R}^{(kn)}(\alpha + \frac{\pi}{2}) \end{aligned}$$

We note that at $\alpha = 0$ the dependence on m_N could be only studied by the investigation of $\sigma_{-1,1}^{(kn)}$ [1], because for the $e_R^- e_L^+$ initial ($-\lambda = \bar{\lambda} = 1$) the dominant channel with the ν_L -exchange is closed and the contribution coming from the diagram with the ν_R -exchange become observable. However, when we take into account the neutrino oscillation angle the situation is totally different. In this case both $\sigma_{-1,1}^{(kn)}$ and $\sigma_{1,1}^{(kn)}$ depend on m_N . The cross sections for initially LL and RR polarized e^-e^+ beams identically vanish and the unpolarized cross section $\sigma^{(kn)}$ is defined by

$$\sigma^{(kn)} = \frac{1}{4} (\sigma_{1,-1}^{(kn)} + \sigma_{-1,1}^{(kn)}) \quad (51)$$

Therefore, we can get the information about m_N investigating the case of the initially unpolarized e^+e^- beams too.

In order to pursue the dependence of the total cross section on α and m_N it is useful to introduce the following quantities

$$\delta_{\lambda\bar{\lambda}}^{(kn)}(\alpha) = \frac{\sigma_{\lambda\bar{\lambda}}^{(kn)}(\alpha) - \sigma_{\lambda\bar{\lambda}}^{(kn)}(0)}{\sqrt{\sigma_{\lambda\bar{\lambda}}^{(kn)}(0)}} \sqrt{LT} \quad (52)$$

$$\delta_{\lambda\bar{\lambda}}(m_N) = \frac{(\sigma_{\lambda\bar{\lambda}})_{SM} - \sigma_{\lambda\bar{\lambda}}^{(11)}}{\sqrt{(\sigma_{\lambda\bar{\lambda}})_{SM}}} \sqrt{LT} \quad (53)$$

where $\sigma_{\lambda\bar{\lambda}}^{(kn)}(\alpha)$ is the ALRM total cross section at the oscillation angle α , $(\sigma_{\lambda\bar{\lambda}})_{SM}$ is the total cross section of the SM and LT is the integrated luminosity. It is obvious that these quantities define the sensitivity of the experiment to the deviations caused by α and m_N and give the deviations expressed in the standard error units.

We start the analysis of the reaction (49) from the case of the $W_1^+ W_1^-$ production. The value of $\delta_{LR}^{(11)}(\alpha)$ is very small over the whole energy region at any m_N . In contrast, for the initial $e_R^- e_L^+$, $\delta_{RL}^{(11)}(\alpha)$ could be large enough to be detectable. In Fig.6 we display $\delta_{RL}^{(11)}(\alpha)$ versus $\sqrt{s}$ for different values of m_N ($g_L = g_R$). We see that $\delta_{RL}^{(11)}(\alpha)$ increases with the growth of m_N and at $m_N = 1$ TeV can reach 1.1σ . The degree of the longitudinal polarization of $e^- e^+$ beams at present is about 0.4. Changing λ ($\bar{\lambda}$) from $-1(1)$ till $-0.4(0.4)$ leads to the near vanishingly small values of $\delta_{\lambda\bar{\lambda}}^{(11)}(\alpha)$. So, the deviations connected with the oscillation angle will be statistically significant at very big m_N ($m_N > 1$ TeV), only.

In Fig.7 we display $\delta_{\lambda\bar{\lambda}}(m_N)$ versus of $\sqrt{s}$ at $m_N = 100$ and 1000 GeV ($g_L = g_R$). From Fig.7 it follows that in the case of the completely RL-polarized $e^- e^+$ -beams we have a good chance to observe the clean signal of the LR symmetry at 2σ level (95% CL) in the LEP II energy region. The increase of the right-handed neutrino mass leads to the growth of $\delta_{\lambda\bar{\lambda}}(m_N)$. It is worth stressing that $\delta_{LR}(m_N)$ reaches its peak at energy of about 190 GeV and then starts to decrease till the Z_2 -resonance.

The investigation $W_1^+ W_2^-$ pair production does not give any information about α . As the analysis shows $\delta_{\lambda\bar{\lambda}}^{(12)}(\alpha)$ is very small at any values of m_N . In Fig.8 we display $\sigma_{1,-1}^{(12)}$ and $\sigma_{-1,1}^{(12)}$ versus $\sqrt{s}$ for $m_N = 100$ and 1000 GeV ($g_L = g_R$). The decrease of $\sigma_{-1,1}^{(12)}$ for a very massive right-handed neutrino is due to the fact that at $s \ll m_N^2$ the functions $\mathcal{D}_{2\nu}^{(kn)}$ and $\mathcal{D}_{3R}^{(kn)}$ decreases with the growth of m_N .

In Fig.9 we depict the total cross section of the reaction $e^+ e^- \rightarrow W_2^+ W_2^-$ as a function of $\sqrt{s}$ for the different values of the right-handed neutrino masses. Again we are limited by consideration of the symmetric case. For the LR polarized $e^+ e^-$ beams the total cross section is smaller on the two orders of magnitude. We see that the decrease of the total cross section for the great values of m_N also takes place. In the case $m_N = 1$ TeV, $\sigma_{LR}^{(22)}$ reaches its maximum in an energy range of about 1.6 TeV and only after that starts to decrease accordingly to (39).

4 Conclusions

Within the asymmetric left-right model (ALRM) $\nu_e - N_e$ and $\nu_e - \nu_\chi$ oscillations have been considered. We have obtained the solution of the Schrödinger equation for a neutrino propagating through matter and magnetic fields. The neutrino behavior within the ALRM and the standard model of electroweak interactions is qualitatively similar but differs if we compare numbers (for the oscillation parameters). In particular the contributions coming from the $SU(2)_R$ sector of the extended model could be sizable. For example, for the convection zone of the sun they could reach 10 percent of its $SU(2)_L$ counterpart. It is possible that nature realized a rather complicated scenario of neutrinos allowing them to have multipole moments, mixings, different helicities and masses. Therefore, some kind of exact solution is useful when investigating the stars interior as well as the neutrino properties. For example, neutrino spin-flavor precession could have remarkable consequences for supernova bursts where $H_\perp \sim 10^{12}$ G and the condition $n_n > n_e$ is

realized in the core region.

Analytical expressions of the cross section for the reaction $e^- e^- \rightarrow W_k^- W_n^-$ in the case of the polarized e^- -beams have been obtained within the framework of the ALRM. The cross section asymptotic behaviour is in accordance with unitarity bounds for initial $e_L^- e_R^-$ beam only. In the case $k = n = 1$ this reaction mainly goes for LR polarized electrons. As the SM predicts σ_{LR} being equal to zero the observation of the nonvanishing cross section will be a clean signal of the LR symmetry of the electroweak interaction. The obtained $\sigma_{LR}^{(11)}$ has maximum in the energy region of about 200 GeV. With the increase of m_N , $\sigma_{LR}^{(11)}$ grows and could reach the values of the order of 0.3 pb at $m_N = 1.4$ TeV.

Concerning $e^- e^- \rightarrow W_1^- W_2^-$, the leading contribution comes from LR-polarized electron beams. Again at the increasing m_N the enhancement mechanism of the total cross section takes place. At the chosen values of SP's the maximum of $\sigma^{(12)}$ is in an energy region ~ 700 GeV. Having reached its maximum the total cross section starts to decrease to the constant values.

The $e^- e^- \rightarrow W_2^- W_2^-$ production is mainly due to RR- polarized electron beams. Its dependence on m_N is very similar to that for $\sigma^{(11)}$ and $\sigma^{(12)}$.

As a final remark, we remind that the general feature of the reactions discussed above is their extreme sensitivity to the such parameters of the ALRM as ξ , λ and g_R .

Turning to $e^+ e^- \rightarrow W_k^+ W_n^-$ reaction, analogous investigation has shown that the dependence on m_N is also the general feature of the total cross section for all its final states while the dependence on α could not be measurable.

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7 Figure captions

- Fig.1** The Feynmann diagrams giving the contributions to the electroweak multipole moments.
- Fig.2** The Feynmann diagrams corresponding to the process $e^-e^- \rightarrow W_k^- W_n^-$.
- Fig.3** The total cross section of the reaction $e_L^- e_R^- \rightarrow W_1^- W_1^-$ versus $\sqrt{s}$ at $m_N = 1$ (dashed line) and $m_N = 1.4 \text{ TeV}$ (solid line).
- Fig.4** $\sigma_{LR}^{(12)}$ as a function of $\sqrt{s}$ at $m_N = 1$ (dashed line) and $m_N = 1.4 \text{ TeV}$ (solid line).
- Fig.5** $\sigma_{LR}^{(22)}$ versus $\sqrt{s}$ at $m_N = 1$ (dashed line) and $m_N = 1.4 \text{ TeV}$ (solid line).
- Fig.6** $\delta_{RL}^{(11)}(\alpha)$ as function of $\sqrt{s}$ at $m_N = 1$ (dashed line) and $m_N = 1.4 \text{ TeV}$ (solid line).
- Fig.7** $\delta_{\lambda\bar{\lambda}}^{(11)}(m_N)$ versus $\sqrt{s}$. The dotted (dash-dotted) line corresponds to the initial $e_L^- e_R^+$ at $m_N = 100 \text{ (1000) GeV}$. The dashed (solid) one does to the initial $e_R^- e_L^+$ at $m_N = 100 \text{ (1000) GeV}$.
- Fig.8** The total cross section of the reaction $e^+e^- \rightarrow W_1^+ W_2^-$ versus $\sqrt{s}$: a) at $m_N = 100 \text{ GeV}$ for $\lambda = -\bar{\lambda} = 1$ - dashed line; b) at $m_N = 100 \text{ (1000) GeV}$ for $-\lambda = \bar{\lambda} = 1$ - solid (dash-dotted) line.
- Fig.9** The total cross section of the reaction $e^-e^+ \rightarrow W_2^- W_2^+$ versus $\sqrt{s}$. The solid and dashed lines correspond to $-\lambda = \bar{\lambda} = 1$ at $m_N = 100$ and 1000 GeV , respectively.

Fig. 1

Fig. 2

Fig. 3

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