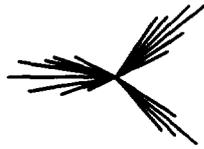




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**Vanishing cosmological constant in elementary
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Vanishing cosmological constant in elementary particles theory

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Abstract

The quest of a vanishing cosmological constant is considered in the simplest anomaly-free chiral gauge extension of the electroweak standard model where the new physics is limited to a well defined additional flavordynamics above the Fermi scale, namely up to a few TeVs by matching the gauge coupling constants at the electroweak scale, and with an extended Higgs structure. In contrast to the electroweak standard model, it is shown how the extended scalar sector of the theory allows a vanishing or a very small cosmological constant. The details of the cancellation mechanism are presented. At accessible energies the theory is indistinguishable from the standard model of elementary particles and it is in agreement with all existing data.

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All astronomical surveys agree that there is no evidence for any spacetime distortion due to a nonvanishing cosmological constant [1] which is many orders of magnitude smaller than that estimated in theories of elementary particles. Up to distances which are accessible to astronomers, about 10 billion light-years, or 10^{28} cm, the magnitude of the cosmological constant must be smaller than $10^{-56} \text{ cm}^{-2} \approx 10^{-84} \text{ GeV}^2$. The possible presence of the cosmological constant Λ in the Einstein's field equations

$$R_{\mu\nu} - \left(\frac{1}{2}R - \Lambda\right) g_{\mu\nu} = \frac{8\pi G}{c^4} T_{\mu\nu} \quad (1)$$

can eventually be vindicated by measuring exponential deviations from the standard matter dominated spatially flat Friedman-Robertson-Walker universe scale factor $R(t) \sim t^{2/3}$. The influence of matter on the metric is determined by the energy-momentum tensor $T_{\mu\nu}$. The component T_{00} is the energy density and the coefficient $c^4/8\pi G \approx 5 \times 10^{45} \text{ g} \times \text{cm} \times \text{sec}^{-2}$, where G is the Newtonian gravitational constant, measures the elasticity of the vacuum. The cosmological constant is, according to astronomical evidence, very close to zero. There is no understanding of why Λ should be close or equal to zero. This question of the cosmological constant problem, which consists in understanding a possible cancellation, is widely regarded as one of the most significant mysteries of the modern cosmology. A proposal for solution, due to Baum-Hawking-Coleman [2], has attracted much attention. It argues that the vanishing of the cosmological constant is closely related to a wormhole-induced quantum instability of the theory. The observable value of the cosmological constant could not be an absolutely fundamental c-number parameter but a dynamical quantum variable with the meaning of topological changes, such as baby universes and wormholes [3]. Nowhere, in physics, we find a greater divergence between theory and experiment than in the cosmological constant problem. As emphasized by Weinberg [4], for a particle physicist, all the values in the observationally allowed range, extending up to values that would make up most of the critical density required in a spatially flat Robertson-Walker universe, seem ridiculously implausible. We should remark that all the standard cosmological model is faced by a severe hierarchical problem, the vacuum energy of the actual universe is extremely fine tuned to

zero [5].

In elementary particle theory the underlying gauge symmetry is larger than that of the actual vacuum whose symmetry is the combination of the color and the abelian electromagnetic factors $G_{3_C 1_{em}} \equiv SU(3)_C \otimes U(1)_{em}$. The full gauge symmetry of the standard model for the nongravitational interactions [6]

$$G_{3_C 2_L 1_Y} \equiv SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$$

is restored at the Fermi scale

$$\Lambda_F = \frac{1}{2^{1/4} \sqrt{G_F}} \approx 246 \text{ GeV}$$

($G_F \approx 1.166 \times 10^{-5} \text{ GeV}^{-2}$) equivalent to a time of order 10^{-11} sec. This scale is set by the decay constant of the three Goldstone bosons transformed via the Higgs-Kibble mechanism [7] into the longitudinal components of the weak gauge bosons. The underlying $G_{3_C 2_L 1_Y}$ symmetry is not manifest in the structure of the vacuum and nature realizes the mechanism of spontaneous symmetry breaking $G_{3_C 2_L 1_Y} \rightarrow G_{3_C 1_{em}}$ where the $G_{3_C 2_L 1_Y}$ symmetry is broken because the associated vacuum state is not invariant anymore. Spontaneous symmetry breaking preserves the renormalizability of the original gauge theory even after symmetry breaking, giving us a renormalizable theory of massive vector bosons [8].

Let us consider a scalar field with a ϕ^4 interaction

$$\begin{aligned} \mathcal{L} &= \partial_\mu \phi \partial^\mu \phi - V(\phi); \\ V(\phi) &= -\frac{1}{2} m^2 \phi^2 + \frac{1}{4} \lambda \phi^4 \end{aligned} \tag{2}$$

containing the discrete symmetry $\phi \rightarrow -\phi$,

$$\mathcal{L}(\phi) \rightarrow \mathcal{L}(-\phi) = \mathcal{L}(\phi).$$

In the electroweak standard model with the multiplet of scalar fields transforming under $G_{3_C 2_L 1_Y}$ as

$$\Phi \equiv \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \sim (\mathbf{1}, \mathbf{2}, +1) \tag{3}$$

the scalar potential can be written in a ϕ^4 fashion

$$V(\Phi^\dagger\Phi) = a\Phi^\dagger\Phi + b(\Phi^\dagger\Phi)^2. \quad (4)$$

After the process of spontaneous symmetry breaking, the neutral component of the scalar doublet gets a vacuum expectation value so that

$$\langle\Phi\rangle = \langle\phi^0\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \sigma \end{pmatrix} \quad (5)$$

and in terms of σ the potential is

$$V(\sigma) = \frac{a}{2}\sigma^2 + \frac{b}{4}\sigma^4.$$

Now we define

$$a \equiv -m^2, \quad b \equiv \lambda; \quad m > 0$$

and the potential becomes

$$V(\sigma) = -\frac{1}{2}m^2\sigma^2 + \frac{1}{4}\lambda\sigma^4 \quad (6)$$

with the minima determined by the conditions

$$V' \equiv \frac{\partial V}{\partial \sigma} = \sigma(-m^2 + \lambda\sigma^2) = 0$$

and

$$V'' \equiv \frac{\partial^2 V}{\partial \sigma^2} = -m^2 + 3\lambda\sigma^2 > 0$$

which show that a particle would rather not sit at the vacuum state $\phi = 0$. Instead, it moves down the potential to a lower-energy state given by the bottom of one of the wells

$$\sigma_{\pm} = \pm \left(\frac{m^2}{\lambda} \right)^{\frac{1}{2}} \quad (7)$$

where the potential takes the value

$$V(\sigma_{\pm}) = -\frac{m^4}{4\lambda} \quad (8)$$

and $V''(\sigma_{\pm}) = 2m^2$ is the curvature of the potential about the true ground state associated to the mass M of the physical boson by

$$M^2 = V''(\sigma_{\pm}) = 2m^2 = 2\lambda\sigma_{\pm}^2. \quad (9)$$

Further, we make the simple choice

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \sigma + H \end{pmatrix} \quad (10)$$

where H is the neutral physical Higgs boson. The potential given in Eq. (4) would become at tree level

$$V(H) = -\frac{m^4}{4\lambda} - m^2 H^2 + \lambda\sigma H^3 + \frac{\lambda}{4} H^4 \quad (11)$$

then, with Eq. (9), we get

$$V(H) = -\frac{m^4}{4\lambda} - \frac{1}{2} M^2 H^2 + \lambda\sigma H^3 + \frac{\lambda}{4} H^4. \quad (12)$$

The field independent constant term $-m^4/4\lambda$ preceding the mass term is the energy density of the vacuum given by the component

$$\langle T_0^0 \rangle \equiv \rho_{\Lambda} \propto V(\sigma_{\pm}) = -\frac{m^4}{4\lambda} \quad (13)$$

of the stress tensor for a scalar field. Following Zel'dovich [5], the contribution of the vacuum energy density plays the role of a cosmological constant

$$\Lambda = \frac{8\pi G}{c^2} \rho_{\Lambda} = \frac{8\pi G}{c^4} V(\sigma_{\pm}) = -\frac{2\pi G}{c^4} m^2 \sigma_{\pm}^2$$

which through the relation involving the Fermi constant and σ

$$\frac{G_F}{\sqrt{2}} = \frac{1}{2\sigma^2}$$

and Eq. (9) yields

$$\Lambda = -\frac{\pi G}{2\sqrt{2}c^4 G_F} M^2 \approx -1.3 \times 10^{-33} M^2 \quad (14)$$

with Λ as a functional dependence of M^2 only. According to the astronomical bound [1]

$$|\Lambda_{\text{obs}}| < 10^{-56} \text{ cm}^{-2} \approx 10^{-84} \text{ GeV}^2,$$

corresponding to

$$|\rho_\Lambda| < 5 \times 10^{-29} \text{ g/cm}^3 \approx 10^{-47} \text{ GeV}^4$$

if σ_\pm are absolute minima of the potential, $V(\sigma_\pm) < V(0)$, and taking the bound on the Higgs boson mass $M > 10 \text{ GeV}$ [9] it follows that

$$|\Lambda/\Lambda_{\text{obs}}| \approx 10^{52}. \tag{15}$$

Therefore in the standard model of elementary particles it is expected a large vacuum energy. The spontaneous symmetry breaking of the gauge symmetry, necessary for generating masses, introduces a vacuum energy into the theory. It is tempting to require $\rho_\Lambda = 0$ which can be accomplished by adding the constant $+m^4/4\lambda$ to the Lagrangian which does not affect the equations of motion or the quantization of the theory. Spontaneous symmetry breaking implies a large cosmological constant in the standard Higgs model. Anyway the supersymmetric standard model as well as superstring theories [10] have this same problem which is the most drastic of all fine-tuning and naturalness problem. A mechanism by which some form of supersymmetry makes ρ_Λ vanish was considered recently by Witten [11]. Also technicolor with extended technicolor, which has spontaneous breaking at least up to 100s of TeV, has a terrible cosmological constant problem [12].

There are no experimental data that directly run against the predictions of the standard theory of elementary particles. All experimental data are consistent with the standard model and so theoretical extensions must be motivated by attempting to understand features that are accommodated in the standard model but not explained by it. In this respect, many theorists believe that there should be some new physics lurking at the TeV scale, accessible to the next generation of colliders, or better yet, at present colliders such as the Tevatron at Fermilab or LEP II at CERN. The simplest chiral gauge extension of the standard model

gauge group that one could consider is

$$G_{3_C 3_L 1_X} \equiv SU(3)_C \otimes SU(3)_L \otimes U(1)_X.$$

Although there exist several models with the $G_{3_C 3_L 1_X}$ gauge symmetry [13] recent proposals have a different representation content and a quite different new physics at no high energy scale [14]. There exist several distinct possibilities depending on the way the electric charge operator for fermions

$$Q/e = \frac{1}{2}(\lambda_3^L + \xi \lambda_8^L) + X \quad (16)$$

is embedded in the neutral generators of the $G_{3_C 3_L 1_X}$ group where ξ is the embedding parameter. Notice that the electric charge operator can be written as

$$Q/e = \frac{1}{2}\lambda_3^L + \frac{1}{2}Y$$

where

$$Y = \xi \lambda_8^L + 2X$$

is the hypercharge operator of the gauge group of the standard model.

The salient features of these models which clarify loose ends of the standard model can be quoted as: 1) In the standard model, each family of fermions is anomaly free. This is true for many extensions including grand unified theories and supersymmetric models. In the $G_{3_C 3_L 1_X}$ each family is anomalous but different families are not replicas of one another, and the anomalies cancel when the number of families are taken into account, and to be a multiple of the number of colors. This novel method of anomaly cancellation requires that at least one family transforms differently from the others, thus breaking generation universality. The $G_{3_C 3_L 1_X}$ is the most economical gauge group which admits such fermion representation and gives the first step for understanding the flavor question [14,15]; 2) The $G_{3_C 3_L 1_X}$ allows to give a natural answer to the family replication question and furthermore gives some indication as to why the *top* quark is so heavy [14,16]; 3) The electroweak mixing angle free parameter, θ_{1V} , of the standard model is limited from above. In the $\xi = -\sqrt{3}$ and $\xi = 1/\sqrt{3}$ cases we

obtain the bounds $\sin^2 \theta_W < 1/4$ and $\sin^2 \theta_W < 3/4$, respectively [17]; 4) Electric charge quantization is another achievement within $G_{3_C 3_L 1_X}$ symmetry. In particular in models with $0, \pm 1$ charged leptons there is always a family transforming as $(\mathbf{1}, \mathbf{3}_L, 0)$ under $G_{3_C 3_L 1_X}$. In this generation there is charge quantization in the fashion of grand unified theories. From anomaly cancellation follows the electric charge quantization in the other families [15]; 5) One of the quark families transforms differently from the other two. Using experimental input on neutral bosons mixing, the third family must be the one that is singled out, at least up to small family mixing [18].

The phenomenology of models with extended gauge symmetries provides the existence of extra gauge bosons and new exotic fermions and their discovery would be a definitive signal of new physics [19]. Let us summarize the principal consequences of $G_{3_C 3_L 1_X}$ models concerning physics beyond the standard model: 1) There are five additional gauge bosons. In the model with the embedding parameter $\xi = -\sqrt{3}$, there are a neutral Z' and four bileptons, (U^{--}, V^-) with lepton number $L = +2$ and (U^{++}, V^+) with $L = -2$. Here $L = L_e + L_\mu + L_\tau$ and there is not conservation of the family lepton number L_i , ($i = e, \mu, \tau$) [20]. At e^-p colliders such as HERA or LEP-II-LHC were considered the prospects of searching for bileptons. At LEP-II-LHC are expected more than 280 events per year provided that the mass of bileptons is less than 1 TeV [21]; 2) An interesting connection constraining $U(1)_{em}$ electromagnetic gauge invariance and the nature of neutrino is realized. In some $G_{3_C 3_L 1_X}$ models the masslessness of the photon prevents the neutrino from acquiring Majorana mass [22]; 3) The $G_{3_C 3_L 1_X} \rightarrow G_{3_C 2_L 1_Y}$ breaking scale is estimated by running $\sin^2 \theta_W$ towards large values which give the upper bound of 1.7 TeV [23] and thus new physics may possibly be in the range accessible to accelerator experiments; 4) Lepton number may be explicitly broken by trilinear scalar self couplings. This leads to neutrino masses proportional to the cube of the corresponding charged lepton mass, with consequences for solar neutrinos and for hot dark matter [24]; 5) The $G_{3_C 3_L 1_X}$ is a phenomenologically viable symmetry for having large magnetic moment for the electron neutrino while keeping its mass naturally small, needed in one proposed solution for the solar neutrino problem [25]; 6) The Yukawa

couplings of the $G_{3_C 3_L 1_X}$ models automatically contain a Peccei-Quinn symmetry. This symmetry can be extended to the entire Lagrangian solving the strong CP problem [26]; 6) Conservation of the leptobarion number $F \equiv B + L$ forbids the existence of massive neutrinos and the neutrinoless double beta decay. Explicit or/and spontaneous breaking of F implies that the neutrinos have an arbitrary mass. The neutrinoless double beta decay also has channels that do not depend explicitly on the neutrino mass [27]; 7) A supersymmetric $G_{3_C 3_L 1_X}$ model has a possible reduction to the standard $G_{3_C 2_L 1_Y}$ model with two doublets at the electroweak energy scale. Because of the existence of cubic invariants in the superpotential of the larger theory, the reduced Higgs potential is not that of the minimal supersymmetric standard model [28]. To explore the Higgs sector at the electroweak energy scale, it is important to realize that even if supersymmetry exists, the minimal supersymmetric standard model is not the only possibility for two Higgs doublets. A first example based on E_6 particle content left-right supersymmetric model has already been discovered [29]; 8) Finally, notice that using the lightest leptons as the particles which determine the approximate symmetry, if each family of fermions is treated separately, the $SU(4)_L$ is the highest symmetry group to be considered in the electroweak sector [30]. Here we find the Ockam razor for direct chiral gauge extensions without exotic charged leptons.

Let us examine the possible cancellation among vacuum contribution terms in the framework of $G_{3_C 3_L 1_X}$ gauge extensions. Consider the simplest case where the scalar fields are attributed as [14,31]

$$\eta = \begin{pmatrix} \eta^0 \\ \eta_1^- \\ \eta_2^+ \end{pmatrix} \sim (\mathbf{1}, \mathbf{3}, 0); \quad \rho = \begin{pmatrix} \rho^+ \\ \rho^0 \\ \rho^{++} \end{pmatrix} \sim (\mathbf{1}, \mathbf{3}, 1); \quad \chi = \begin{pmatrix} \chi^- \\ \chi^{--} \\ \chi^0 \end{pmatrix} \sim (\mathbf{1}, \mathbf{3}, -1), \quad (17)$$

which give the following pattern of symmetry breaking

$$SU(3)_L \otimes U(1)_N \xrightarrow{\langle \chi \rangle} SU(2)_L \otimes U(1)_Y \xrightarrow{\langle \rho, \eta \rangle} U(1)_{\text{em}}.$$

The neutral components of the Higgs triplets of the Eqs. (17) develop the vacuum expectation values v_η , v_ρ and v_χ , respectively, with $(246 \text{ GeV})^2 \equiv v_W^2 = v_\eta^2 + v_\rho^2$. Since the

fields are worked as perturbations around the stable vacuum we define

$$\varphi = v_\varphi + \xi_\varphi + i\zeta_\varphi, \quad (18)$$

$\varphi = \eta^0, \rho^0, \chi^0$. With the three triplets of the Eqs. (17) we can write the more general, renormalizable and gauge invariant Higgs potential

$$\begin{aligned} V_T(\eta, \rho, \chi) = & \mu_1^2 \eta^\dagger \eta + \mu_2^2 \rho^\dagger \rho + \mu_3^2 \chi^\dagger \chi + \lambda_1 (\eta^\dagger \eta)^2 + \lambda_2 (\rho^\dagger \rho)^2 + \lambda_3 (\chi^\dagger \chi)^2 + \\ & + (\eta^\dagger \eta) \left[\lambda_4 (\rho^\dagger \rho) + \lambda_5 (\chi^\dagger \chi) \right] + \lambda_6 (\rho^\dagger \rho) (\chi^\dagger \chi) + \lambda_7 (\rho^\dagger \eta) (\eta^\dagger \rho) + \\ & + \lambda_8 (\chi^\dagger \eta) (\eta^\dagger \chi) + \lambda_9 (\rho^\dagger \chi) (\chi^\dagger \rho) + \left(\frac{f_1}{2} \varepsilon^{ijk} \eta_i \rho_j \chi_k + \text{H. c.} \right), \end{aligned} \quad (19)$$

where the μ 's, λ 's and f_1 are coupling constants. The leptobarionic number $F = L + B$ is conserved in Eq.(19), where L and B are the total leptonic and barionic numbers, respectively [27]. In order to avoid linear terms in φ fields the potential shifted according to Eq. (18) requires the validity of the relations

$$\mu_1^2 + 2\lambda_1 v_\eta^2 + \lambda_4 v_\rho^2 + \lambda_5 v_\chi^2 + f_1 \frac{v_\rho v_\chi}{2v_\eta} = 0, \quad (20a)$$

$$\mu_2^2 + 2\lambda_2 v_\rho^2 + \lambda_4 v_\eta^2 + \lambda_6 v_\chi^2 + f_1 \frac{v_\eta v_\chi}{2v_\rho} = 0, \quad (20b)$$

$$\mu_3^2 + 2\lambda_3 v_\chi^2 + \lambda_5 v_\eta^2 + \lambda_6 v_\rho^2 + f_1 \frac{v_\eta v_\rho}{2v_\chi} = 0. \quad (20c)$$

where $\text{Im} f_1 = 0$. The scalar sector of the $G_{3_c 3_L 1_X}$ model which we are considering here was carried out in Ref. [32] in the approximation $v_\chi \approx -f_1 \gg v_\eta, v_\rho$. This approximation leads to the conditions

$$\lambda_4 \approx 2 \frac{\lambda_2 v_\rho^2 - \lambda_1 v_\eta^2}{v_\eta^2 - v_\rho^2}, \quad \lambda_5 v_\eta^2 + 2\lambda_6 v_\rho^2 \approx -\frac{v_\eta v_\rho}{2}. \quad (21)$$

Since we are interested in the $G_{3_c 3_L 1_X}$ vacuum contribution to the cosmological constant let us replace Eq. (18) in Eq. (19) and use the constraints (20) and (21). Eliminating terms representing vacuum fluctuations and maintaining only terms in lower order in $1/v_\chi$ we obtain, after to require the vanishing of the remaining potential terms,

$$v_\chi = v_\eta \left[\frac{\lambda_1 (4v_\eta^2 - 3v_W^2)}{\lambda_3 (v_W^2 - 2v_\eta^2)} \right]^{1/4} \quad (22)$$

where $\lambda_3 < 0$ in order to obtain real square masses for the neutral scalar bosons [32].

The model of the Refs. [14] requires the sextet

$$S = \begin{pmatrix} \sigma_1^0 & s_2^+ & s_1^- \\ s_2^+ & S_1^{++} & \sigma_2^0 \\ s_1^- & \sigma_2^0 & S_2^{--} \end{pmatrix} \sim (\mathbf{1}, \mathbf{6}^*, 0) \quad (23)$$

of scalar fields besides the triplets of the Eqs. (17). Here we are considering $\langle \sigma_1 \rangle = 0$ for maintaining zero neutrino mass [27]. Thus, the Eqs. (18) have only one more component $\sigma_2 = v_2 + \xi_2 + i\zeta_2$, where $\langle \sigma_2 \rangle = v_2$. In this case we get the Higgs potential as

$$\begin{aligned} V_S(\eta, \rho, \chi, S) = & V_T + \mu_4^2 \text{Tr}(S^\dagger S) + \lambda_{10} \text{Tr}^2(S^\dagger S) + \lambda_{11} \text{Tr}(S^\dagger S)^2 + \\ & + [\lambda_{12} (\eta^\dagger \eta) + \lambda_{13} (\rho^\dagger \rho) + \lambda_{14} (\chi^\dagger \chi)] \text{Tr}(S^\dagger S) + \\ & + \left(\frac{f_2}{2} \rho_i \chi_j S^{ij} + \text{H.c.}\right), \end{aligned} \quad (24)$$

where V_T is given in Eq.(19). The constraint equations equivalent to Eqs. (20) are expressed now by

$$\mu_1^2 + 2\lambda_1 v_\eta^2 + \lambda_4 v_\rho^2 + \lambda_5 v_\chi^2 + 2\lambda_{12} v_2^2 + f_1 \frac{v_\eta v_\chi}{2v_\eta} = 0, \quad (25a)$$

$$\mu_2^2 + 2\lambda_2 v_\rho^2 + \lambda_4 v_\eta^2 + \lambda_6 v_\chi^2 + 2\lambda_{13} v_2^2 + f_1 \frac{v_\eta v_\chi}{2v_\rho} + f_2 \frac{v_\chi v_2}{2v_\rho} = 0, \quad (25b)$$

$$\mu_3^2 + 2\lambda_3 v_\chi^2 + \lambda_5 v_\eta^2 + \lambda_6 v_\rho^2 + 2\lambda_{14} v_2^2 + f_1 \frac{v_\eta v_\rho}{2v_\chi} + f_2 \frac{v_\rho v_2}{2v_\chi} = 0, \quad (25c)$$

$$\mu_4^2 + 2(2\lambda_{10} + \lambda_{11})v_2^2 + \lambda_{12} v_\eta^2 + \lambda_{13} v_\rho^2 + \lambda_{14} v_\chi^2 + f_2 \frac{v_\rho v_\chi}{4v_2} = 0. \quad (25d)$$

The approximation

$$v_\chi \approx -f_1 \approx |f_2| \gg v_\eta, v_\rho, v_2, \quad (26)$$

with $v_W^2 = v_\eta^2 + v_\rho^2 + v_2^2$ leads to [32]

$$4(2\lambda_{10} + \lambda_{11}) \approx \lambda_1 \approx \lambda_2 \approx \lambda_{12} \approx \lambda_{13} \approx \frac{\lambda_4}{2}. \quad (27)$$

Following analogous procedure as in the three triplets case, we replace Eqs. (26) and (27) into the potential V_S , Eq. (24), and we maintain only the vacuum expectation value contributions. Next, if we require that the potential terms vanish we obtain

$$(-\lambda_6 + \lambda_9) v_W^2 + (-1 - \lambda_5 + \lambda_6 + \lambda_8) v_\eta^2 - \lambda_3 v_\chi^2 + (-3\lambda_{14} + \lambda_6 - \lambda_9) v_2^2 = 0. \quad (28)$$

This is the condition for an arbitrarily small effective cosmological constant in the case where the symmetric sextet of scalar fields play an essential role in order to generate the charged lepton masses.

We have presented a way for bringing the energy density of the vacuum to zero, associated with the flat empty world of Minkowski, in a gauge model of elementary particles. In all elementary particle theories, including supersymmetry and grand unification, it is predicted a very large value. Specially, we find remarkable that a vanishing cosmological constant can be obtained in a relatively simple extension of the standard model which has several replies for some fundamental questions that plague the physics of elementary particles. However, even if it is shown that the value of Λ can be zero, the cosmological constant problem still remains, whether Λ actually does vanish exactly and identically due to some symmetry or another unknown deep physical connection.

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