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Connecting Dark Energy to Neutrinos with an Observable Higgs Triplet

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Abstract

To connect the scalar field (acceleron) responsible for dark energy to neutrinos, the usual strategy is to add unnaturally light neutral singlet fermions (right-handed neutrinos) to the Standard Model. A better choice is actually a Higgs triplet, through the coupling of the acceleron to the trilinear Higgs triplet-double-doublet interaction. This hypothesis predicts an easily observable doubly-charged Higgs boson at the forthcoming Large Hadron Collider (LHC).

The existence of dark energy [1] may be attributed to a scalar field called the acceleron (or quintessence) [2] whose equation of motion involves a term of negative pressure, allowing the present Universe to expand at an accelerated rate. The acceleron may also form a condensate and couple to matter in such a way that the observed neutrino masses are dynamical quantities. This is the scenario of mass varying neutrinos [3], motivated by the proximity of the effective mass scale of dark energy to that of neutrinos, which may have some interesting consequences [4, 5].

To make the connection, the usual strategy is to introduce 3 right-handed neutrinos N_i, i.e. 3 neutral fermion singlets under the electroweak $SU(2)_L \times U(1)_Y$ gauge group. However, contrary to the cherished expectation that m_{N_i} should be very large (thereby triggering the canonical seesaw mechanism [6] and yielding naturally small Majorana neutrino masses m_{ν_i}), they have to be very small, i.e. of order eV, to be compatible with dark energy. In view of this problem, alternative mechanisms for the origin of m_{ν_i} should be explored [7].

In the Standard Model, naturally small Majorana neutrino masses come from the unique dimension-five operator [8]

$$\mathcal{L}_{eff} = \frac{f_{ij}}{\Lambda}(\nu_i \phi^0 - l_i \phi^+)(\nu_j \phi^0 - l_j \phi^+) + H.c., \tag{1}$$

which can be realized at tree level in exactly 3 ways [9], one of which is of course the canonical seesaw mechanism with 3 right-handed neutrinos. Another way is to add a Higgs triplet [10]

$$\Delta = \begin{pmatrix} \xi^+/\sqrt{2} & \xi^{++} \\ -\xi^0 & -\xi^+/\sqrt{2} \end{pmatrix} \tag{2}$$

with trilinear couplings to both the lepton doublets (ν_i, l_i) and the Higgs doublet $\Phi = (\phi^+, \phi^0)$, i.e.

$$\mathcal{L}_{int} = f_{ij}[\nu_i \nu_j \xi^0 + \frac{1}{\sqrt{2}}(\nu_i l_j + l_i \nu_j)\xi^+ + l_i l_j \xi^{++}] + \mu \Phi^\dagger \Delta \tilde{\Phi} + H.c., \tag{3}$$

where $\tilde{\Phi} = (\bar{\phi}^0, -\phi^-)$. As a result [11],

$$(\mathcal{M}_\nu)_{ij} = \frac{2 f_{ij} \mu \langle \phi^0 \rangle^2}{m^2_{\xi^0}}. \quad (4)$$

If $\mu = \mu(\mathcal{A})$, i.e. a function of the acceleron field $\mathcal{A}$, then this is in fact a natural realization of mass varying neutrinos with m_ξ of order the electroweak scale.

In all previous proposals of neutrino mass with a Higgs triplet, there is no compelling reason for m_ξ to be this low. One possible exception [12] is the case of large extra space dimensions, where m_ξ should be below whatever the cutoff energy scale is, but that is only a phenomenological lower bound. On the other hand, if dark energy is indeed connected to neutrinos through the Higgs triplet, then at least ξ^{++} will be unambiguously observable at the forthcoming Large Hadron Collider (LHC).

Consider the most general Higgs potential consisting of Φ and Δ, i.e.

$$\begin{aligned} V &= m^2(\Phi^\dagger \Phi) + M^2(Tr\Delta^\dagger \Delta) + \frac{1}{2}\lambda_1(\Phi^\dagger \Phi)^2 + \frac{1}{2}\lambda_2(Tr\Delta^\dagger \Delta)^2 \\ &+ \frac{1}{2}\lambda_3(Tr\Delta^\dagger \Delta \Delta^\dagger \Delta) + \lambda_4(\Phi^\dagger \Phi)(Tr\Delta^\dagger \Delta) + \lambda_5(\Phi^\dagger \Delta^\dagger \Delta \Phi) \\ &+ \mu(\tilde{\Phi}^\dagger \Delta^\dagger \Phi + \Phi^\dagger \Delta \tilde{\Phi}). \end{aligned} \quad (5)$$

Let $\langle \phi^0 \rangle = v$ and $\langle \xi^0 \rangle = u$, then

$$v[m^2 + \lambda_1 v^2 + \lambda_4 u^2 - 2\mu u] = 0, \quad (6)$$

$$u[M^2 + (\lambda_2 + \lambda_3)u^2 + \lambda_4 v^2] - \mu v^2 = 0. \quad (7)$$

For $|\mu| << |m|, |M|$, and

$$m^2 < 0, \quad \lambda_1 M^2 - \lambda_4 m^2 > 0, \quad (8)$$

we have the unique solution

$$v^2 \simeq -\frac{m^2}{\lambda_1}, \quad u \simeq \frac{\mu v^2}{M^2 + \lambda_4 v^2}. \quad (9)$$

The Higgs triplet masses are then

$$m_{\xi^{++}}^2 \simeq M^2 + (\lambda_4 + \lambda_5)v^2, \tag{10}$$

$$m_{\xi^{+}}^2 \simeq M^2 + (\lambda_4 + \lambda_5/2)v^2, \tag{11}$$

$$m_{\xi^{0}}^2 \simeq M^2 + \lambda_4 v^2. \tag{12}$$

Once produced, the decay of ξ^{++} into two charged leptons is an unmistakable signature with negligible background. Its decay branching fractions also map out $|f_{ij}|$, i.e. the entire neutrino mass matrix up to an overall scale [12].

In a model of neutrino dark energy (νDE), the neutrino mass m_ν is a dynamical quantity. It is assumed to be a function of a scalar field $\mathcal{A}$ (the acceleron) with a canonically normalized kinetic term and $\partial m_\nu / \partial \mathcal{A} \neq 0$. In the nonrelativistic limit, m_ν depends on the total density n_ν of the thermal background of neutrinos and antineutrinos, and the energy or effective potential of the system is given by

$$V = m_\nu n_\nu + V_0(m_\nu). \tag{13}$$

The thermal background and the scalar potential $V_0(m_\nu)$ will act in opposite directions and at any instant of time, the minimum of the effective potential is given by

$$V'(m_\nu) = n_\nu + V_0'(m_\nu) = 0. \tag{14}$$

We assume the curvature scale of V to be much larger than the Hubble expansion rate, so that the adiabatic approximation is valid. In other words, the solution of Eq. (14) for m_ν is assumed to be valid instantaneously.

For an adiabatic expansion of the Universe, the density of matter varies with the scale factor as

$$\rho \propto R^{-3(1+\omega)}, \tag{15}$$

where ω is a time-independent parameter, which enters in the following simple equation of state:

$$\rho(t) = \omega p(t). \tag{16}$$

In a νDE model, it was shown that ω satisfies the equation

$$\omega + 1 = -\frac{V'(m_\nu)}{m_\nu V} = \frac{\Omega_\nu}{\Omega_\nu + \Omega_{DE}}, \tag{17}$$

where $\Omega_{DE} = \rho_{DE}/\rho_c$ is the contribution of $V_0(m_\nu)$ to the energy density and $\Omega_\nu = n_\nu/\rho_c$ is the neutrino energy density. Since the observed value [1]

$$\omega = -0.98 \pm 0.12$$

is close to -1 at the present time, Ω_ν should be much less compared to Ω_{DE}. These considerations restrict the possibilities of the form of the potential. For small $d\omega/dn_\nu$, the variable mass of the neutrino is proportional to the neutrino density to the power ω:

$$m_\nu \propto n_\nu^\omega.$$

The above general considerations are valid, independent of the details of the particular model of neutrino mass. However, most phenomenological implications are specific to such details, with a few general features which are common to all models [4].

In the present scenario, for the effective neutrino mass to vary, we have to associate the acceleron field $\mathcal{A}$ with the trilinear coupling of Δ with Φ, so that the effective neutrino mass becomes dependent on the field $\mathcal{A}$. This simply means that we set $\mu = \mu(\mathcal{A})$ in the scalar potential of Eq. (5). As for the self-interactions of $\mathcal{A}$, we may assume for example the following potential:

$$V_0 \ = \ \Lambda^4 \ \log(1 + |\bar{\mu}/\mu(\mathcal{A})|). \tag{18}$$

Using Eq. (4), the effective low-energy Lagrangian is then given by

$$-\mathcal{L}_{eff} = f_{ij} \ |\mu(\mathcal{A})| \ \frac{\langle\phi^0\rangle^2}{m_{\xi^0}^2} \ \nu_i \nu_j + H.c. + \Lambda^4 \log(1 + |\bar{\mu}/\mu(\mathcal{A})|), \tag{19}$$

and Eq. (13) is of the form

$$V(x) = ax + b\log\left(1+\frac{c}{x}\right), \tag{20}$$

where $x = m_\nu \propto |\mu(\mathcal{A})|$ and a, b, c are all positive. For $4b/ac << 1$, $x_{min} \simeq b/a$, so

$$m_\nu \propto n_\nu^{-1}, \tag{21}$$

as desired. As a condition of naturalness, it has been argued that the mass of the scalar field should not be larger than the order of 1 eV and $\Lambda \sim 10^{-3}$ eV. In the canonical realization of mass varying neutrinos using right-handed neutrinos N, this would imply small NN Majorana masses as well as tiny νN Dirac masses, which are clearly rather unnatural. Here, the requirement is simply that m_{ξ^0} be of order $\langle\phi^0\rangle$, which is a much more reasonable condition.

Thus the mass of ξ^0 is predicted to be in the range of $80 - 500$ GeV. The lower limit is the present experimental bound from the direct search of the triplet Higgs scalar, while the upper limit comes from the requirement that it should not be too large compared to the electroweak breaking scale, otherwise it would be difficult to explain neutrino masses much below 1 eV. The form of $\mu(\mathcal{A})$ was discussed in the original paper [3] to be $\mu(\mathcal{A}) \sim \lambda\mathcal{A}$ or $\mu(\mathcal{A}) \sim \mu e^{\mathcal{A}^2/f^2}$. We shall not go into the details of this discussion on the dynamics of this model, although some of the generic problems of mass varying neutrinos are common to the present model as well [14].

Depending on the form of $\mu(\mathcal{A})$, global lepton number may be broken spontaneously in such a model of νDE, thereby creating a massless Goldstone boson, i.e. the Majoron. However, as shown below, its coupling to ordinary matter is highly suppressed, hence its existence is acceptable phenomenologically. If we take the case $\mu(\mathcal{A}) \sim \lambda\mathcal{A}$ (where $\mathcal{A}$ is complex), we can express the field $\mathcal{A}$ as

$$\mathcal{A} = \frac{1}{\sqrt{2}}(\rho + \sqrt{2}z)e^{i\varphi}$$

where z is the vacuum expectation value or condensate of $\mathcal{A}$. Similarly,

$$\phi^0 = \frac{1}{\sqrt{2}}(H + \sqrt{2}v)e^{i\theta}, \quad \xi^0 = \frac{1}{\sqrt{2}}(\zeta + \sqrt{2}u)e^{i\eta}, \tag{22}$$

with v and u as the vacuum expectation values of ϕ^0 and ξ^0 respectively. The longitudinal component of the Z boson (G^0), the physical Majoron (J^0) and the massive combination (Ω^0) of $(z\varphi, u\eta, v\theta)$ are given by:

$$\begin{aligned} G^0 &= \frac{v^2\theta + 2u^2\eta}{\sqrt{v^2 + 4u^2}}, \\ J^0 &= \frac{(v^2 + 4u^2)z^2\varphi + v^2u^2\eta - 2u^2v^2\theta}{\sqrt{z^2(v^2 + 4u^2)^2 + u^2v^4 + 4v^2u^4}} \\ \Omega^0 &= \frac{\varphi - \eta + 2\theta}{\sqrt{z^{-2} + u^{-2} + 4v^{-2}}}, \end{aligned} \tag{23}$$

respectively. The heavy Ω^0 is almost degenerate in mass with ζ. They are essentially the reincarnations of ξ^0. The massless J^0 is potentially a problem phenomenologically but its couplings to all leptons are strongly suppressed by $(u/v)^2$, and can safely be neglected in all present experiments.

Since the triplet Higgs scalars cannot be much heavier than the usual Higgs doublet, they should be observable at the LHC as well as the proposed future International Linear Collider (ILC). The phenomenology of such triplet Higgs scalars has been discussed in [12]. The same-sign dileptons will be the most dominating decay modes of the ξ^{++}. Complementary measurements of $|f_{ij}|$ at the ILC by the process $e^-e^-(\mu^-\mu^-) \to l_i^- l_j^-$ would allow us to study the structure of the neutrino mass matrix in detail. Of course, these features are generic to any model with a Higgs triplet as the origin of Majorana neutrino masses. The difference here is that it is also accompanied by the unusual predictions of mass varying neutrinos in neutrino oscillations [4, 15].

In conclusion, we have pointed out in this paper that if the neutrino mass m_ν is dynamical and related to dark energy through the acceleron $\mathcal{A}$, then the most natural mechanism for generating m_ν is that of the Higgs triplet, rather than the canonically assumed right-handed neutrino. The mass scale of the triplet Higgs scalars is predicted to be close to that of electroweak symmetry breaking, hence it has an excellent chance of being observed at the LHC and ILC. Aspects of this model relating to cosmology and neutrino oscillations are similar to other existing models of dark energy.

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