Naturally light neutrinos in \textit{Diracon} model

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Abstract

We propose a simple model for Dirac neutrinos where the smallness of neutrino mass follows from a parameter \( \kappa \) whose absence enhances the symmetry of the theory. Symmetry breaking is performed in a two-doublet Higgs sector supplemented by a gauge singlet scalar, realizing an accidental global U(1) symmetry. Its spontaneous breaking at the few TeV scale leads to a physical Nambu–Goldstone boson – the \textit{Diracon}, denoted \( D \) – which is restricted by astrophysics and induces invisible Higgs decays such as \( h \to DD \). The scheme provides a rich, yet very simple scenario for symmetry breaking studies at colliders such as the LHC.

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I. INTRODUCTION

Establishing whether neutrinos are their own anti-particles continues to challenge experimentalists [1,2]. Likewise, the mechanism responsible for generating small neutrino masses remains as elusive as ever. It is well-known that, in gauge theories, the detection of neutrinoless double beta decay would signify that neutrinos are of Majorana type [3,4]. Although experimental confirmation of the Majorana hypothesis may come in the not too distant future [5], so far the possibility remains that neutrinos can be Dirac particles, despite the fact that the general theoretical expectation is that they are Majorana fermions [6] as given, for example, in Weinberg’s dimension five operator [7]. Moreover, the most widely studied mechanism to account for the smallness of neutrino masses relative to the charged fermion masses invokes their Majorana nature, namely, the conventional high-scale type-I [6,8–11] or type-II [6,11,12] seesaw mechanism. The same happens in low-scale variants of the seesaw mechanism, see [13] for a review.

Accommodating the possibility of naturally light Dirac neutrinos constitutes a double challenge. One approach is to supplement the standard SU(3)_c ⊗ SU(2)_L ⊗ U(1)_Y electroweak gauge structure by using extra flavor symmetries implying a conserved lepton number, so as to obtain Dirac neutrinos, as suggested in [14,15]. Another approach would be to appeal to the existence of extra dimensions, such as in warped scenarios [16]. Alternatively, one may extend the gauge group itself, for example, by using the extended SU(3)_c ⊗ SU(3)_L ⊗ U(1)_X gauge structure because of its special features [17]. In this case one can obtain both the lightness as well as the Dirac nature of neutrinos as an outcome [18].

In this letter we focus on the possibility of having naturally light Dirac neutrinos with seesaw-induced masses within the framework of the simplest four-dimensional SU(3)_c ⊗ SU(2)_L ⊗ U(1)_Y gauge structure, without non-Abelian discrete flavor symmetries. To this end we impose a cyclic flavor-blind Z_5 ⊗ Z_3 symmetry in a theory with enlarged symmetry breaking sector : two Higgs doublets and a singlet, see Table I. We find that the resulting model has an accidental spontaneously broken U(1) symmetry that leads to the seesaw mechanism as well as the Dirac nature of neutrinos. The smallness of neutrino mass follows from the smallness of a parameter κ whose absence would increase the symmetry of the electroweak breaking sector, ensuring naturalness in ’t Hooft’s sense. We discuss some phenomenological features of the scheme which follow from the existence of a Diracon, namely, the Nambu–Goldstone boson associated to the spontaneous breaking of the global accidental U(1) symmetry in the scalar sector.
II. THE MODEL

The lepton and scalar boson assignments of the model are summarized in Table I, where a cyclic $Z_5 \otimes Z_3$ symmetry is assumed, so that $\omega^5 = 1$ and $\alpha^3 = 1$. The Abelian $Z_5$ symmetry is used to have Dirac neutrinos in the presence of the additional doublet $\Phi$, forbidding the terms $L \nu_R \tilde{H}$, $\bar{L} \ell_R \bar{\Phi}$, $\nu_R \nu_R$ and $\nu_R \nu_R \sigma$. We have used two Higgs doublets

$$H = \begin{pmatrix} h^+ \\ H^0 \end{pmatrix}, \quad \Phi = \begin{pmatrix} \Phi^0 \\ \phi^- \end{pmatrix},$$

with their conjugates defined as usual, $\tilde{H} = i \sigma_2 H^*$ and $\tilde{\Phi} = i \sigma_2 \Phi^*$. On the other hand, the complementary $Z_3$ symmetry [19, 20] ensures lepton number conservation also at the non-renormalizable level, ruling out possible operators of the type $\nu_R \nu_R \sigma^3$, $\nu_R \nu_R \sigma^8$, etc.

| $SU(2)_L$ | $\bar{L}$ | $\ell_R$ | $\nu_R$ | $H$ | $\Phi$ | $\sigma$ |
|-----------|-----------|----------|---------|-----|-------|--------|
| $Z_5$     | $\omega$  | $\omega^4$ | $\omega$ | $1$ | $\omega^3$ | $\omega$ |
| $Z_3$     | $\alpha^2$ | $\alpha$  | $\alpha$ | $1$ | $1$    | $1$    |

Table I: Lepton and scalar boson assignments of the model, with $\omega^5 = 1$ and $\alpha^3 = 1$.

The gauge– and $Z_5 \otimes Z_3$–invariant Yukawa Lagrangean for the lepton sector turns out to be, symbolically,

$$\mathcal{L}_Y = y^e \overline{L} e_R H + y^\nu \overline{L} \nu_R \Phi + h.c.$$  \hspace{1cm} (1)

where the first term is the standard one responsible for the charged lepton masses, while the second is the one that appears in Fig. 1. As illustrated in the figure, the latter induces nonzero neutrino masses

$$m_\nu = \kappa y^\nu v_\nu^2 v_H m_\Phi^2$$  \hspace{1cm} (2)

where we denote the three vacuum expectation values as $v_\sigma \equiv \langle \sigma \rangle$, $v_\Phi \equiv \langle \Phi \rangle$, $v_H \equiv \langle H \rangle$, and $\kappa$ is a dimensionless parameter in the scalar potential. For simplicity we have omitted generation indices. Notice that the smallness of neutrino mass depends not only on the Yukawa coupling $y^\nu$ but is also related to the smallness of $\kappa$, very much like in the type-II–like seesaw mechanism.
III. SCALAR SECTOR

The $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ invariant scalar potential consistent with the global $Z_5 \otimes Z_3$ symmetry is given by

$$V = -\mu_H^2 H^\dagger H - \mu_H^2 \Phi^\dagger \Phi - \mu^2 \sigma^\dagger \sigma + \lambda_H H^\dagger H H^\dagger H + \lambda_\Phi \Phi^\dagger \Phi^\dagger \Phi + \lambda_\sigma \sigma^\dagger \sigma^\dagger \sigma$$

$$+ \lambda_{H\Phi} H^\dagger H \Phi^\dagger \Phi + \lambda_{\Phi H} \Phi^\dagger \Phi H + \lambda_{\sigma H \sigma} \sigma^\dagger H^\dagger H + \lambda_{\Phi \Phi \sigma} \sigma \Phi^\dagger \Phi$$

$$+ \kappa \left( \hat{H}^\dagger \Phi \sigma^2 + h.c. \right)$$

where, after acquiring vacuum expectation values (vevs) the fields are shifted as follows,

$$H^0 = \frac{1}{\sqrt{2}} (v_H + R_H + iI_H), \quad \sigma = \frac{1}{\sqrt{2}} (v_\sigma + R_\sigma + iI_\sigma) \quad \text{and} \quad \Phi^0 = \frac{1}{\sqrt{2}} (v_\Phi + R_\Phi + iI_\Phi) \tag{4}$$

so the extremum conditions are,

$$\mu_H^2 = \frac{1}{2} \left( 2\lambda_H v_H^2 + \lambda_\sigma H v_\sigma^2 + \lambda_{H\Phi} v_\Phi^2 - \frac{\kappa v_\Phi^2 v_\sigma}{v_H} \right) \tag{5}$$

$$\mu_\Phi^2 = \frac{1}{2} \left( \lambda_\Phi v_\Phi^2 + \lambda_\sigma v_\sigma^2 + 2\lambda_\Phi v_\Phi^2 - \frac{\kappa v_H v_\sigma^2}{v_\Phi} \right)$$

and from these one can derive a “seesaw–type relation” amongst the vacuum expectation values given as,

$$v_\Phi \approx \kappa v_H \left( \frac{1}{\lambda_\Phi v_\Phi^2 + \lambda_\sigma v_\sigma^2 - \frac{\kappa v_\Phi^2}{v_H}} \right). \tag{6}$$

Notice that $v_\sigma \neq 0$ is required in order to drive $v_\Phi \neq 0$, see Fig. [1]. Moreover one sees that the smallness of $v_\Phi$ is directly related to the smallness of $\kappa$. In other words, $\Phi$ acquires an “induced” vev $v_\Phi$ whose smallness is associated to the symmetry enhancement that results from the absence of $\kappa$. In this limit there would be a second $U(1)$ symmetry whose

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1 This scalar potential is shared by other models with Majorana neutrinos. See for example Ref. [21].
associated Nambu-Golstone boson is the pseudoscalar $A$, see below. This means that the “induced” vev $v_{\Phi}$ is always very much suppressed w.r.t. the standard $v_H$, responsible for generating the W boson mass. In short the model has a double vev hierarchy

$$v_{\sigma} \gtrsim v_H \gg v_{\Phi}.$$  \hspace{1cm} (7)

The two hierarchies are consistent with the minimization of the potential. The first is a mild hierarchy, ensuring adequate couplings for the Diracron, while the second one is a strong yet “natural” hierarchy, because it is related to the enhanced symmetry which would result from the absence of $\kappa$ in the Lagrangean, even though, in practice, it can not be strictly realized, since we need $v_{\Phi} \neq 0$ for a realistic scheme.

With the information above one can immediately work our the Higgs mass spectrum. Out of the ten scalars, eight from the two-doublet structure, plus two from the extra complex singlet, we are left with seven physical ones after projecting out the three longitudinal degrees of freedom of the massive $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge bosons. These correspond to three physical CP even scalars, one of which is the 125 GeV state discovered at the LHC [22–24], two physical CP odd scalars, and one electrically charged scalar. The mass squared matrices for the CP-even and CP-odd sectors, in the weak basis $(H, \sigma, \Phi)$ are given below,

$$M_R^2 = \begin{pmatrix}
2\lambda_H v_H^2 + \frac{\kappa v_{\Phi}^2 v_H}{2v_H} & (\lambda_{\sigma H} v_H - \kappa v_{\Phi}) v_{\sigma} & \lambda_{\sigma H} v_H v_{\Phi} - \frac{\kappa v_{\Phi}^2}{2} \\
(\lambda_{\sigma H} v_H - \kappa v_{\Phi}) v_{\sigma} & 2\lambda_\sigma v_{\sigma}^2 & (\lambda_{\sigma \Phi} v_{\Phi} - \kappa v_H) v_{\sigma} \\
\lambda_{\sigma H} v_H v_{\Phi} - \frac{\kappa v_{\Phi}^2}{2} & (\lambda_{\sigma \Phi} v_{\Phi} - \kappa v_H) v_{\sigma} & 2\lambda_\Phi v_{\Phi}^2 + \frac{\kappa v_H v_{\Phi}^2}{2v_{\Phi}}
\end{pmatrix}$$  \hspace{1cm} (8)

and

$$M_I^2 = \kappa \begin{pmatrix}
\frac{v_{\Phi}^2 v_H}{2v_H} & v_{\sigma} v_{\Phi} & \frac{v_{\Phi}^2}{2} \\
v_{\sigma} v_{\Phi} & 2v_H v_{\Phi} & v_H v_{\sigma} \\
\frac{v_{\Phi}^2}{2} & v_H v_{\sigma} & \frac{v_H v_{\Phi}^2}{2v_{\Phi}}
\end{pmatrix},$$  \hspace{1cm} (9)

where \(\text{diag}(m_{H_1}^2, m_{H_2}^2, m_{H_3}^2) = O_R M_R^2 O_R^T\) and \(\text{diag}(0, 0, m_A^2) = O_I M_I^2 O_I^T\).

Let us first consider the CP-odd scalars sector, its diagonalization matrix is given by,

$$O_I = \begin{pmatrix}
-\alpha v_H & 0 & \alpha v_{\Phi} \\
-2\alpha v_H v_{\Phi}^2 & \frac{\beta_{\sigma \Phi}}{\alpha} & -2\alpha v_H v_{\Phi}^2 \\
\beta v_{\sigma} v_{\Phi} & 2\beta v_H v_{\Phi} & \beta v_H v_{\sigma}
\end{pmatrix}$$  \hspace{1cm} (10)

where

$$\alpha = \frac{1}{\sqrt{v_H^2 + v_{\Phi}^2}} \quad \text{and} \quad \beta = \frac{1}{\sqrt{v_H^2 (v_{\sigma}^2 + 4v_{\Phi}^2) + v_{\Phi}^2 v_{\Phi}^2}}.$$  \hspace{1cm} (11)
Hence one finds that the mass–eigenstate profiles are given by

\[ G^0 = \alpha (-v_H I_H + v_\Phi I_\Phi) \]
\[ D = \alpha \beta (-2v_H v_\Phi^2 I_H + \frac{v_\sigma}{\alpha^2} I_\sigma - 2v_H^2 v_\Phi I_\Phi) \]
\[ A = \beta (v_\sigma v_\Phi I_H + 2v_H v_\Phi I_\sigma + v_H v_\sigma I_\Phi). \] (12)

One sees that the projective nature of Eq. (9) (two–dimensional null space) clearly implies two massless states whose profiles follow just from symmetry reasons. Due to the smallness of \( v_\Phi \), the main components of \( G^0, D \) and \( A \) are the imaginary parts of the \( SU(2)_L \) Higgs doublet \( H \), the singlet \( \sigma \) and the doublet \( \Phi \), respectively. Indeed the first massless CP–odd eigenvector \( G^0 \) pointing mainly along \( H \) corresponds to the unphysical longitudinal mode of the Z boson. The second massless state \( D \) is mainly singlet and we call it the Diracon, i.e. the physical Nambu-Golstone boson associated to the accidental U(1) symmetry. It is the analogue of the Majoron present in the “123” type-II seesaw scheme of [11], and is associated with the type-II Dirac neutrino seesaw mechanism illustrated in Fig. 1. On the other hand the massive pseudoscalar state \( A \) pointing mainly along the weak isodoublet direction has mass

\[ m_A^2 = \frac{\kappa (v_H^2 (v_\sigma^2 + 4v_\Phi^2) + v_\sigma^2 v_\Phi^2)}{2v_H v_\Phi}, \] (13)

which would vanish in the (unphysical) limit \( \kappa = v_\sigma = v_\Phi = 0 \).

Turning now to the charged sector we have, in the basis \((h^\pm, \phi^\pm)\), the following mass squared matrix,

\[ M_{H^\pm}^2 = \begin{pmatrix}
\lambda_H v_H^2 + \frac{\kappa v_\sigma^2}{v_H} & \lambda_H v_H v_\Phi - \kappa v_\sigma^2 \\
-\lambda_H v_H v_\Phi - \kappa v_\sigma^2 & \lambda_H v_H + \frac{\kappa v_\sigma^2}{v_H}
\end{pmatrix} \] (14)

whose eigenstates are the longitudinal W–boson and a physical state \( H^\pm \) of (squared) mass

\[ m_{H^\pm}^2 = (v_H^2 + v_\Phi^2) \left( \lambda_H + \frac{\kappa v_\sigma^2}{v_H v_\Phi} \right). \] (15)

Notice that, taking into account the smallness of the neutrino mass, i.e. \( \kappa \ll 1 \), and using Eq. 6 one finds that the Higgs mass spectrum further simplifies to,

\[ M_R^2 \approx \begin{pmatrix}
2\lambda_H v_H^2 & \lambda_H v_H v_\sigma & 0 \\
\lambda_H v_H v_\sigma & 2\lambda_\sigma v_\sigma^2 & 0 \\
0 & 0 & \lambda_H v_H^2 + \lambda_\sigma v_\sigma^2
\end{pmatrix} \] (16)

\[ m_A^2 \approx \frac{\lambda_H v_H^2 + \lambda_\sigma v_\sigma^2}{2} \] (17)

\[ m_{H^\pm}^2 \approx \frac{\lambda_H v_H^2}{2} + (\lambda_H + \lambda_H^\prime) \frac{v_H^2}{2}. \] (18)
Comparing the CP–even and CP–odd sectors it follows that \( m_{H_3} \approx m_A \). Hence by using Eq. (18) and Eq. (17) we find that the following mass relation holds,

\[
m_{H^\pm}^2 - m_A^2 \approx \lambda' \frac{v_H^2}{2}.
\]

IV. PHENOMENOLOGICAL CONSIDERATIONS

The above model of electroweak breaking is rather similar to the inert doublet model \[25\], implying the absence of tree-level flavor-changing neutral currents. There are, however, important new features. A noticeable difference of this model when compared with various variants of two–Higgs doublet models is the existence of the accidental U(1) symmetry. This global symmetry is spontaneously broken by the vev of \( \sigma \) implying the existence of a corresponding Nambu–Goldstone boson given in Eq. (12). Its couplings to neutrinos can be easily obtained from Noether’s theorem. Likewise one can determine its coupling to charged leptons, for instance electrons. The latter would lead to excessive stellar cooling through the Compton–like process \( \gamma + e \to e + D \) \[26\], unless

\[
|g_{eeD}| = \left| (\mathcal{O}_I)_{21} \frac{m_e}{v_H} \right| \lesssim 10^{-13}
\]

hence, using Eq. (10), one finds

\[
2\alpha\beta v_\Phi^2 \lesssim \frac{10^{-13}}{m_e}
\]

where \( \alpha \) and \( \beta \) are given in Eq. (11). Taking into account that \( v_H = \sqrt{v_{SM}^2 - v_\Phi^2} \) (where \( v_{SM} = 246 \text{ GeV} \)), one writes Eq. (20) only in terms of \( v_\sigma \) and \( v_\Phi \). The allowed region of these vevs is delimited by the bound on \( g_{eeD} \) as illustrated in Fig. 2. The shaded area is the region allowed by stellar energy loss limits.

Figure 2: The shaded region is allowed by stellar energy loss limits.
Let us now comment on boundedness conditions and vacuum stability. Taking \( \kappa \ll 1 \), one can see that the copositivity criterium applies so one can easily obey the relevant conditions \[27\]. Concerning stability and perturbativity, we find extented regions of consistency as a result of the presence of the extra scalar boson states \[28, 29\]. Moreover, in the limit \( m_{H_3} \sim m_A \sim m_{H^\pm} \), it is well–known that the oblique S,T,U parameters are well under control, so that these precision observables does not pose great problems either \[30, 31\].

Concerning collider phenomenology, notice that the physical Nambu–Goldstone boson induces invisible Higgs boson decays \( h \rightarrow DD \). These decays are rather analogous to the invisible CP–even Higgs decays into Majorons which are present in Majorana neutrino schemes, such as the “123” seesaw \[11\] with spontaneous lepton number violation \[32\]. Likewise, one has the new pseudoscalar decays \( A \rightarrow hD \) and \( A \rightarrow DDD \) in addition to the Standard Model decay modes. For charged scalars, there are also new decay channels into leptons, i.e. \( H^\pm \rightarrow \ell^\pm \nu_R \), which should be taken into account in search analyses \[33\]. Future experimental searches at the LHC should probe the theory in a rather significant way within a relatively wide region of parameters.

V. SUMMARY AND CONCLUSIONS

We presented a very simple model where neutrinos are Dirac fermions and their mass can naturally arise from TeV–scale physics, associated to a small parameter \( \kappa \) whose absence would enhance the symmetry of the theory. This is realized in an enlarged scalar sector consisting of two doublets and a singlet Higgs carrying an accidentally conserved global U(1) charge. Its spontaneous violation leads to the existence of a physical Nambu–Goldstone boson which is restricted by astrophysics. Let us mention that the presence of gravity could induce masses for neutrinos \[34\] and/or the Dirac \[35\] breaking the Abelian discrete symmetries. This breaking may, however, be avoided if the latter are part of a gauge discrete symmetry \[36\]. We have discussed the symmetry structure of the model, the connection to neutrino mass generation, and indicated how it provides new collider signatures induced by invisible Higgs boson decays. In summary, the model provides an interesting scheme for neutrino mass generation. Its scalar sector constitutes an interesting alternative for electroweak symmetry breaking studies, both theoretically and experimentally. Its simplicity, its close connection to neutrino masses and the presence of an accidental global U(1) symmetry give it unique features. Details as well as additional phenomenological features will be discussed elsewhere.

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[1] A. S. Barabash, Phys. Atom. Nucl. **74**, 603 (2011), 1104.2714.
[2] F. T. Avignone, S. R. Elliott, and J. Engel, Rev. Mod. Phys. **80**, 481 (2008), 0708.1033.
[3] J. Schechter and J. Valle, Phys.Rev. **D25**, 2951 (1982).
[4] M. Duerr, M. Lindner, and A. Merle, JHEP **1106**, 091 (2011), 1105.0901.
[5] A. Gando et al. (KamLAND-Zen collaboration) (2016), 1605.02889.
[6] J. Schechter and J. Valle, Phys.Rev. **D22**, 2227 (1980).
[7] S. Weinberg, Phys. Rev. **D22**, 1694 (1980).
[8] M. Gell-Mann, P. Ramond, and R. Slansky (1979), print-80-0576 (CERN).
[9] T. Yanagida (KEK lectures, 1979), ed. O. Sawada and A. Sugamoto (KEK, 1979).
[10] R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 91 (1980).
[11] J. Schechter and J. W. F. Valle, Phys. Rev. **D25**, 774 (1982).
[12] G. Lazarides, Q. Shafi, and C. Wetterich, Nucl. Phys. **B181**, 287 (1981).
[13] S. M. Boucenna, S. Morisi, and J. W. Valle, Adv. High Energy Phys. **2014**, 831598 (2014).
[14] A. Aranda et al., Phys. Rev. **D89**, 033001 (2014), 1307.3553.
[15] S. Kanemura, K. Sakurai, and H. Sugiyama, Phys. Lett. **B758**, 465 (2016), 1603.08679.
[16] P. Chen et al., JHEP **01**, 007 (2016), 1509.06683.
[17] M. Singer, J. Valle, and J. Schechter, Phys.Rev. **D22**, 738 (1980).
[18] J. W. F. Valle and C. A. Vaquera-Araujo, Phys. Lett. **B755**, 363 (2016), 1601.05237.
[19] E. Ma and R. Srivastava, Phys. Lett. **B741**, 217 (2015), 1411.5042.
[20] E. Ma, N. Pollard, R. Srivastava, and M. Zakeri, Phys. Lett. **B750**, 135 (2015), 1507.03943.
[21] W. Wang and Z.-L. Han (2016), 1605.00239.
[22] G. Aad et al. (ATLAS Collaboration), Phys.Lett. **B716**, 1 (2012), 1207.7214.
[23] S. Chatrchyan et al. (CMS Collaboration), Phys.Lett. **B716**, 30 (2012), 1207.7235.
[24] G. Aad et al. (ATLAS, CMS), Phys. Rev. Lett. **114**, 191803 (2015), 1503.07589.
[25] G. C. Branco, P. Ferreira, L. Lavoura, M. Rebelo, M. Sher, and J. P. Silva, Phys. Reports **516**, 1 (2012).
[26] N. Viaux, M. Catelan, P. B. Stetson, G. Raffelt, J. Redondo, A. A. R. Valcarce, and A. Weiss, Phys. Rev. Lett. **111**, 231301 (2013), 1311.1669.
[27] K. Kannike, Eur. Phys. J. **C72**, 2093 (2012), 1205.3781.
[28] C. Bonilla, R. M. Fonseca, and J. W. F. Valle, Phys. Lett. **B756**, 345 (2016), 1506.04031.
[29] C. Bonilla, R. M. Fonseca, and J. W. F. Valle, Phys. Rev. **D92**, 075028 (2015), 1508.02323.
[30] W. Grimus, L. Lavoura, O. Ogreid, and P. Osland, J.Phys. **G35**, 075001 (2008), 0711.4022.
[31] W. Grimus, L. Lavoura, O. Ogreid, and P. Osland, Nucl.Phys. **B801**, 81 (2008), 0802.4353.
[32] M. A. Diaz et al., Phys. Rev. **D58**, 057702 (1998), hep-ph/9712487.
[33] G. Aad et al. (ATLAS), JHEP 03, 127 (2016), 1512.03704.
[34] G. Dvali and L. Funcke, Phys. Rev. D93, 113002 (2016), 1602.03191.
[35] S. R. Coleman, Nucl. Phys. B310, 643 (1988).
[36] L. M. Krauss and F. Wilczek, Phys. Rev. Lett. 62, 1221 (1989).