Invisible K_L decays as a probe of new physics

S. N. Gninenko[1] and N. V. Krasnikov[1,2]

[1]*Institute for Nuclear Research of the Russian Academy of Sciences, 117312 Moscow, Russia*
[2]*Joint Institute for Nuclear Research, 141980 Dubna, Russia*

(Received 10 March 2015; published 10 August 2015)

The decay $K_L \to invisible$ has never been experimentally tested. In the Standard Model (SM), its branching ratio for the decay into two neutrinos is helicity suppressed and predicted to be $\mathrm{Br}(K_L \to \nu\bar{\nu}) \lesssim 10^{-10}$. We consider several natural extensions of the SM, such as two-Higgs-doublet (2HDM), 2HDM and light scalar, and mirror dark matter models, whose main feature is that they allow us to avoid the helicity suppression factor and lead to an enhanced $\mathrm{Br}(K_L \to invisible)$. For the decay $K_L \to \nu\bar{\nu}$, the smallness of the neutrino mass in the considered 2HDM model is explained by the smallness of the second Higgs doublet vacuum expectation value. The small nonzero value of the second Higgs isodoublet can arise as a consequence of nonzero quark condensate. We show that taking into account the most stringent constraints from the $K \to \pi + invisible$ decay, this process could be in the region of $\mathrm{Br}(K_L \to invisible) \simeq 10^{-8}$–$10^{-6}$, which is experimentally accessible. In some scenarios, the $K_L \to invisible$ decay could still be allowed while the $K \to \pi + invisible$ decay is forbidden. The results obtained show that the $K_L \to invisible$ decay is a clean probe of new physics scales well above 100 TeV that is complementary to rare $K \to \pi + invisible$ decay, and they provide a strong motivation for its sensitive search in a near-future experiment.

DOI: 10.1103/PhysRevD.92.034009 PACS numbers: 13.20.Eb, 12.60.Fr, 13.35.Hb, 14.80.-j

I. INTRODUCTION

In the Standard Model (SM), the branching ratios of the $K^+ \to \pi^+ + invisible$ and $K_L \to \pi^0 + invisible$ decays are predicted to be [1]

$$\mathrm{Br}(K_L \to \pi^0\nu\bar{\nu}) = (2.6 \pm 0.4) \times 10^{-11}, \tag{1}$$

$$\mathrm{Br}(K^+ \to \pi^+\nu\bar{\nu}) = (8.5 \pm 0.7) \times 10^{-11}, \tag{2}$$

with the invisible final state represented by neutrino pairs. A strong comparison between experiment and theory is possible due to the accuracy of both the measurements and the SM calculations of these observables. A discrepancy would signal the presence of physics beyond the Standard Model (BSM), making the precision measurements of these decays an effective probe to search for it, see e.g. Refs. [1–6].

The branching ratio of the $K_L \to invisible$ decay in the SM is predicted to be very small compared to those of Eqs. (1) and (2) for ν masses lying in the sub-eV region favored by observations of ν oscillations [7]. Indeed, the K_L has zero spin, and it cannot decay into two massless neutrinos, as it contradicts momentum and angular-momentum conservation simultaneously. For the case of massive ν's, their spins in the K_L rest frame must be opposite and, therefore, one of them is forced to have the "wrong" helicity. This results in the $K_L \to \nu\bar{\nu}$ decay rate being proportional to the ν mass squared $\Gamma(K_L \to \nu\bar{\nu}) \propto (\frac{m_\nu}{m_{K_L}})^2 \lesssim 10^{-17}$, assuming $m_\nu \lesssim 1$ eV. However, if one takes the direct experimental upper limit on the ν_τ mass $m_{\nu_\tau} < 18.2$ MeV [7], the predicted branching ratio, calculated at the quantum loop level, is [8]

$$\mathrm{Br}(K_L \to \nu\bar{\nu}) \simeq 10^{-10}. \tag{3}$$

Therefore, an observed $\mathrm{Br}(K_L \to invisible) \gg 10^{-10}$ would unambiguously signal the presence of BSM physics.

The decay $K_L \to invisible$ has never been experimentally tested. Since long ago, it was recognized that this decay "would be interesting to explore, but its detection looks essentially impossible. New ingenious experimental ideas are required" [8]. Recently, an approach for performing such experiments by using the $K^+n \to K^0p$ (or $K^-p \to \bar{K}^0n$) charge-exchange reaction as a source of well-tagged K^0's has been reported [9]. At the same time, the first experimental bound, $\mathrm{Br}(K_L \to invisible) \lesssim 6.3 \times 10^{-4}$, has been set from existing experimental data. It has been shown that compared to this limit, the expected sensitivity of the proposed search is at least 2 orders of magnitude higher—$\mathrm{Br}(K_L \to invisible) \lesssim 10^{-6}$ per $\simeq 10^{12}$ incident kaons. It could be further improved by utilizing a more careful design of the experiment, thus making the region $\mathrm{Br}(K_L \to invisible) \simeq 10^{-8}$–$10^{-6}$, or even below, experimentally accessible [9].

Being motivated by these considerations, we discuss in this work several natural extensions of the SM and show that taking into account the most stringent constraints from the measured $K^+ \to \pi^+ + invisible$ decay rate, the decay $K_L \to invisible$ could occur at the level $\mathrm{Br}(K_L \to invisible) \simeq 10^{-8}$–$10^{-6}$. The main feature of the considered models, that leads to an enhanced branching ratio for $K_L \to invisible$, compared to $K^+ \to \pi^+ + invisible$, is that they allow us to avoid the helicity suppression factor $(\frac{m_\nu}{m_{K_L}})^2$ in the SM, while profiting from its larger phase space due to the decay into two light weakly interacting particles. In addition, there might be the case in which $K_L \to invisible$ could still be kinematically allowed, while $K^+ \to \pi^+ + invisible$ is forbidden. Additional motivation to search for the K_L (and K_S) invisible decay is related to precision tests of the $K^0 - \bar{K}^0$ system by using the Bell-Steinberger unitarity relation [9]. This relation connects CP and CPT violation in the mass matrix to CP and CPT violation in all decay channels of neutral kaons and is a powerful tool for testing CPT invariance with neutral kaons [10]. The question of how much the invisible decays of K_S or K_L can influence the precision of the Bell-Steinberger analysis still remains open [11]. All this makes the future searches for this decay mode very interesting and complementary to the study of the $K \to \pi + invisible$ decays.

II. $K_L \to \nu\bar{\nu}$ DECAY IN MODEL WITH ADDITIONAL SCALAR DOUBLET

Consider now the $K_L \to \nu\bar{\nu}$ decay in *the two-Higgs-doublet model* (2HDM) with an additional heavy Higgs doublet H_2. This type of 2HDM model can introduce flavor-changing neutral currents, and provide explanations of the origin of dark matter and CP violation; see e.g. Ref. [12]. The interaction of the heavy isodoublet field H_2 with quarks, leptons, and the standard Higgs isodoublet H leading to the $K_L \to \nu\bar{\nu}$ decay has the form

$$L_{\rm int} = h_{2\tau}\bar{L}_\tau \tilde{H}_2 \nu_{\tau_R} + h_{2d_L s_R}\bar{Q}_{1L} H_2 s_R + \delta m^2_{HH_2} H^+ H_2 + \text{H.c.} - M^2_{H_2} H_2^+ H_2, \qquad (4)$$

where $L_\tau = (\nu_{\tau_L}, \tau_L)$, $Q_{1L} = (u_L, d_L)$, $H_2 = (H_2^+, H_2^0)$, $\tilde{H}_2 = ((H_2^0)^*, -(H_2^+)^*)$ and $h_{2\tau}, h_{2d_L s_R}$ are Yukawa coupling constants. Note that in general the second Higgs isodoublet H_2 will have nonzero Yukawa interactions with other quark and lepton fields, but since we are interested mainly in the $K_L \to \nu\bar{\nu}$ decay, we have written explicitly only the Yukawa interactions important for us. In the considered model, the neutrinos acquire nonzero Dirac masses $m_{\nu_\tau} = h_{2\tau}\langle H_2 \rangle$ due to the nonzero vacuum expectation value of the second Higgs isodoublet $\langle H_2 \rangle = \frac{\delta m^2_{HH_2}}{M^2_{H_2}} \langle H \rangle$ ($\langle H \rangle = 174$ GeV), and the smallness of the Dirac neutrino masses is a consequence of the $\langle H_2 \rangle$ smallness. The smallness of $\langle H_2 \rangle$ is due to the assumed large value of M_{H_2} or (and) the small value of $\delta m^2_{HH_2}$ [13]. For instance, for $m_{\nu_\tau} = 0.1$ eV, $h_{2\tau} = 0.1$ and $M_{H_2} = 10^5$ GeV, we find that $\frac{\delta m^2_{HH_2}}{M^2_{H_2}} = 0.6 \times 10^{-11}$ and $\delta m^2_{HH_2} = 0.06$ GeV^2. It is interesting to note that for $\delta m^2_{HH_2} = 0$ the $\langle H_2 \rangle = 0$ at classical level, but the spontaneous symmetry breaking of $SU_L(3) \otimes SU_R(3)$ chiral symmetry in QCD leads to nonzero vacuum expectation values for the Higgs fields [15]. Really, for the nonzero Yukawa interaction $L_{H_2 Q_1 d} = h_{2d_L d_R}\bar{Q}_{1L} H_2 d_R + \text{H.c.}$, due to the nonzero vacuum expectation value of the quark condensate $\langle \bar{d}d \rangle = -\frac{f_\pi^2 m_\pi^2}{(m_u + m_d)}$ ($f_\pi = 93$ MeV), the field $\langle H_2 \rangle$ acquires a nonzero vacuum expectation value $\langle H_2 \rangle = \frac{\langle \bar{d}d \rangle}{2h_{2d_L d_R} M^2_{H_2}}$. Numerically, for $h_{2\tau} = h_{2d_L d_R} = 1$ and $m_{\nu_\tau} = 0.1$ eV, we find that $M_{H_2} \sim O(10^4)$ GeV. So in this model with $\delta m^2_{HH_2} = 0$, the vacuum expectation value $\langle H_2 \rangle = 0$ at tree level, but the nonzero quark condensate leads to the appearance of a small vacuum expectation value $\langle H_2 \rangle \neq 0$ for the second Higgs isodoublet that explains the smallness of the neutrino masses.

For the case of nonzero neutrino Majorana mass $m_{\nu_{\tau_R}}$, we assume that the mass $m_{\nu_{\tau_R}}$ is small, so the decay $K_L \to \nu_\tau \bar{\nu}_\tau$ is kinematically allowed. Again, as in the previous case, we assume that the Dirac neutrino mass arises due to a nonzero $\langle H_2 \rangle$ vacuum expectation value, and the smallness of the seesaw $m_{\nu_{\tau_R}} = \frac{m^2_{D\nu\tau}}{m_{\nu_{\tau_R}}}$ neutrino mass is again explained due to the smallness of $\langle H_2 \rangle$. The Lagrangian (4) contains $\Delta S = 1$ neutral flavor-changing terms, but for a heavy doublet H_2 it is not dangerous. The effective four-fermion Lagrangian describing the decay $K_L \to \nu_\tau \bar{\nu}_\tau$ has the form

$$L_{\rm eff} = \frac{1}{M_X^2}\bar{d}_L s_R \bar{\nu}_{\tau_L} \nu_{\tau_R} + \text{H.c.}, \qquad (5)$$

where

$$\frac{1}{M_X^2} = \frac{h_{2d_L s_R} h_{2\tau}}{M^2_{H_2}}. \qquad (6)$$

As has been mentioned before, we assume the existence of a small Dirac or Majorana neutrino mass ν_τ. The decay rate of the invisible decay $K_L \to \nu_\tau \bar{\nu}_\tau$ is determined by the formula

$$\Gamma(K_L \to \nu_{L\tau}\bar{\nu}_{R\tau}, \nu_{R\tau}\bar{\nu}_{L\tau}) = \frac{M^5_{K_L}}{16\pi M_X^4}\left(\frac{F_K}{2(m_d + m_s)}\right)^2 K(m_\nu^2/M^2_{K_L}), \qquad (7)$$

where $K(x) = (1 - 4x)^{1/2}$ for a Dirac neutrino with a mass m_{ν_τ} and $K(x) = (1 - x)^2$ for a Majorana neutrino ν_{τ_R} with a mass $m_{\nu_{\tau_R}}$. Here $F_K \approx 160$ MeV is the kaon decay constant and m_s, m_d are the masses of s and d quarks [16]. For $\text{Br}(K_L \to \nu_\tau \bar{\nu}_\tau) = 10^{-6}$, we can test the value of M_X up to [17]

$$M_X \lesssim 0.6 \times 10^5 \text{ GeV} \quad (8)$$

for small Dirac or Majorana neutrino masses $m_{\nu_\tau} \ll M_{K_L}$.

It should be noted that the existence of $\Delta S = 1$ neutral flavor-changing interaction (5) leads to additional contribution to rare decays $K_L \to \pi^0 \nu\bar{\nu}$ and $K^+ \to \pi^+ \nu\bar{\nu}$. The current experimental values are [18,19]

$$\text{Br}(K_L \to \pi^0 \nu\bar{\nu}) < 2.6 \times 10^{-8}, \quad (9)$$

$$\text{Br}(K^+ \to \pi^+ \nu\bar{\nu}) = (17.3^{+11.5}_{-10.5}) \times 10^{-11}, \quad (10)$$

with the SM predictions of (1) and (2), respectively. The measured value (10) for the $\text{Br}(K^+ \to \pi^+ \nu\bar{\nu})$ allows us to set more stringent constraints. Therefore, we restrict ourselves to the calculation of the BSM contribution only to this decay channel by using the effective Lagrangian (5). This leads to the following formula for the differential $K^+ \to \pi^+ \nu\bar{\nu}$ decay width:

$$\begin{aligned} &\frac{d\Gamma^{\text{BSM}}(K^+ \to \pi^+ \nu\bar{\nu})}{dq^2} \\ &= \frac{1}{(2\pi)^3} \cdot \frac{1}{32 M_{K^+}^3} \cdot \frac{(q^2 - m_{\nu_{\tau,R}}^2)^2}{q^2 M_X^4} \\ &\cdot \sqrt{(M_{K^+}^2 + M_{\pi^+}^2 - q^2)^2 - 4 M_{K^+}^2 M_{\pi^+}^2} \\ &\cdot \left[\frac{f_0(q^2)(M_{K^+}^2 - m_{\pi^+}^2)}{2(-m_d + m_s)}\right]^2. \end{aligned} \quad (11)$$

The form factor $f_0(q^2)$ is determined in the standard way as

$$\begin{aligned} \langle \pi | \bar{d}\gamma_\mu s | K \rangle &= f_+(q^2)(P_K + P_\pi)^\mu + f_-(q^2)(P_K - P_\pi)^\mu \\ &= f_+(q^2)\left[(P_K + P_\pi)^\mu - \frac{M_K^2 - M_\pi^2}{q^2} q^\mu\right] \\ &\quad + f_0(q^2)\frac{M_K^2 - M_\pi^2}{q^2} q^\mu, \end{aligned} \quad (12)$$

where $q^\mu = (P_K - P_\pi)^\mu$ and $m_{\nu_R}^2 \le (M_{K^+} - M_{\pi^+})^2$. The form factors f_+ and f_0 are related to the exchange of 1^- and 0^+, respectively. The following relation holds:

$$\begin{aligned} &f_+(0) = f_0(0), \\ &f_0(q^2) = f_+(q^2) + \frac{q^2}{M_K^2 - M_\pi^2} f_-(q^2). \end{aligned} \quad (13)$$

In our calculations we use standard linear parametrization for the form factor $f_0(q^2)$, namely

$$f_0(q^2) = f_0(0)\left(1 + \lambda_0 \frac{q^2}{M_{\pi^+}^2}\right). \quad (14)$$

Numerically, we take $f_0(0) = 0.96$ [20] and $\lambda_0 = -0.06$ [21].

It is convenient to represent the result in terms of the ratio $\beta^{-1} \equiv \frac{\text{Br}(K_L \to \nu\bar{\nu})}{\text{Br}(K^+ \to \pi^+ \nu\bar{\nu})}$, because the ratio β does not depend on the unknown value of M_X. Also, β does not depend on the values of quark masses m_d, m_s. For the case of a massless neutrino, we find that

$$\beta \approx 2 \times 10^{-3}. \quad (15)$$

Note that the smallness of the β is mainly due to the three-body phase space smallness in comparison with two-body phase space. From the difference between the theoretical and experimental values (2) and (10), respectively, by summing up errors of (10) in quadrature we find that the BSM contribution to the $\text{Br}(K^+ \to \pi^+ \nu\bar{\nu})$ is less than

$$\text{Br}^{\text{BSM}}(K^+ \to \pi^+ \nu\bar{\nu}) \lesssim 2.1 \times 10^{-10}. \quad (16)$$

From the limit (16) and the estimate (15), we find that for massless neutrinos

$$\text{Br}(K_L \to \nu\bar{\nu}) \lesssim 10^{-7}. \quad (17)$$

The estimates (15), (17) are valid for a small $m_{\nu_R} \ll M_{\pi^+}$ Majorana mass of the right-handed neutrino. For higher m_{ν_R} values, the limit (17) is more weak; and for the case $M_{K_L} \ge m_{\nu_R} \ge M_{K^+} - M_{\pi^+}$, when the decay $K^+ \to \pi^+ \nu_{\tau_L} \bar{\nu}_{\tau_R}$ is kinematically prohibited, but the decay $K_L \to \nu\bar{\nu}$ is still allowed, the restriction from $K^+ \to \pi^+ \nu\bar{\nu}$ decay does not work.

The measured $(K_L - K_S)$ mass difference strongly restricts [22] the effective $\Delta S = 2$ interaction

$$L_{\bar{s}d\bar{s}d} = \frac{1}{\Lambda_{\bar{s}d\bar{s}d}^2} \bar{s_R} d_L \bar{s_R} d_L + \text{H.c.} \quad (18)$$

Namely [22],

$$\Lambda_{\bar{s}d\bar{s}d} \ge 1.8 \times 10^7 \text{ GeV}. \quad (19)$$

For the model (2) with the additional Higgs doublet $H_2 = (H_2^+, H_{2,1}^0 + iH_{2,2}^0)$, we find that

$$\frac{1}{\Lambda_{\bar{s}d\bar{s}d}^2} = |h_{2d_L s_R}|^2 \left|\frac{1}{M_{H_{2,1}^0}^2} - \frac{1}{M_{H_{2,2}^0}^2}\right| \sim \frac{|h_{2d_L s_R}|^2}{M_{H_2}^2} \cdot \frac{\delta m_{HH_2}^2}{M_{H_2}^2}. \quad (20)$$

Using the bound (19), we can restrict the parameter $\delta m_{HH_2}^2$. For instance, for $M_{H_2} = 10^5$ GeV, $h_{2d_L s_R} = 1$ we find $\delta m_{HH_2}^2 \le 0.3 \times 10^6$ GeV^2, which is much more weak than the estimate of $\delta m_{HH_2}^2$ coming from the neutrino mass.

In the general case, we can have an additional flavor-changing Yukawa interaction $h_{2s_L d_R} \bar{Q}_{2L} H_2 d_R + \text{H.c.}$ $(Q_{2L} = (c_L, s_L)$ in the Lagrangian (4) that leads to the tree-level flavor-changing $\Delta S = 2$ effective interaction $L_{\rm eff} = \frac{h_{2d_L s_R} h^*_{2s_L d_R}}{M^2_{H2}} (\bar{d}_L s_R \bar{d}_R s_L + {\rm H.c.})$. We can simultaneously avoid the $\Delta S = 2$ bound $\Lambda_{\Delta s=2} \equiv (h_{2d_L s_R} h^*_{2s_L d_R})^{-1/2} \cdot M_{H_2} > 1.8 \times 10^7$ GeV and obtain phenomenologically interesting values for $\mathrm{Br}(K_L \to \nu\bar{\nu})$ for small quark Yukawa coupling constants $h_{2d_L s_R}, h_{2s_L d_R}$, a relatively light second Higgs doublet and not a small lepton Yukawa coupling constant $h_{2\tau}$. For instance, for $h_{2d_L s_R} = h_{2s_L d_R} = (1/300)^2$, $h_{2\tau} = 1$ and $M_{H_2} = 300$ GeV, we find that $\Lambda_{\Delta s=2} = 2.7 \times 10^7$ GeV and $\mathrm{Br}(K_L \to \nu\bar{\nu}) = 0.4 \times 10^{-6}$. The existence of a relatively light second Higgs doublet with a mass $M_{H_2} = 300$ GeV does not contradict the LHC data. The best way to look for the second Higgs isodoublet at the LHC is the use of the reaction $pp \to Z^*/\gamma^* \to H_2^+ H_2^- \to \tau^+\tau^-\nu\bar{\nu}$. So the signature is two τ leptons plus nonzero $E^T_{\rm miss}$ in a final state that coincides with the signature used for the search for the direct production of stau leptons at the LHC.

III. $K_L \to \phi\phi$ DECAY IN MODEL WITH ADDITIONAL SCALAR DOUBLET AND SCALAR SINGLET ϕ

Consider now the $K_L \to \phi\phi$ decay in the extension of the SM *with a heavy Higgs doublet H_2 and light neutral scalar singlet field ϕ*. The Yukawa interaction of the heavy isodoublet H_2 with quarks and the interaction of the ϕ field with Higgs isodoublets H_2 and H (Higgs isodoublet of the SM) has the form

$$L_I = h_{2d_L s_R} \bar{Q}_{1L} s_R H_2 + \lambda (H_2^+ H)\phi^2 + \delta m^2_{HH_2} H^+ H_2 + {\rm H.c.} - M^2_{H_2} H_2^+ H_2, \quad (21)$$

where $Q_{1L} = (u_L, d_L)$, $H_2 = (H_2^+, H_2^0)$ and $h_{2d_L s_R}$, λ are Yukawa and Higgs couplings. After electroweak $SU_L(2) \otimes U(1)$ symmetry breaking, a trilinear term describing the transition $H_2 \to \phi\phi$,

$$L_{H_2\phi\phi} = \lambda \langle H \rangle H_2^+ \phi^2 + {\rm H.c.}, \quad (22)$$

arises. The effective Lagrangian

$$L_{\rm eff} = \frac{1}{M_X} \bar{d}_L s_R \phi^2 + {\rm H.c.}, \quad (23)$$

$$\frac{1}{M_X} = \frac{h_{2d_L s_R} \lambda \langle H \rangle}{M^2_{H_2}} \quad (24)$$

describes invisible decay $K_L \to \phi\phi$. Here we assume that the mass of ϕ is less than $M_{K_L}/2$. The decay rate of the invisible decay $K_L \to \phi\phi$ is determined by the formula

$$\Gamma(K_L \to \phi\phi) = \frac{M^3_{K_L}}{8\pi M^2_X} \left(\frac{F_K}{2(m_d + m_s)} \right)^2 K(m^2_\phi / M^2_{K_L}), \quad (25)$$

where $K(x) = (1 - 4x)^{1/2}$. For $\mathrm{Br}(K_L \to \phi\phi) = 10^{-6}$ and $m_\phi \ll M_{K_L}$, we can test the value of M_X up to

$$M_X \lesssim 10^{10}\ {\rm GeV}. \quad (26)$$

For $\lambda = 1$ and $h_{2d_L s_R} = 1$, the mass of the second Higgs isodoublet can be tested up to $M_{H_2} \leq 10^6$ GeV.

The bound (16) allows us to restrict the $K_L \to \phi\phi$ decay in full analogy with the previous model. Namely, in the model with the effective Lagrangian (22), the $K_L \to \phi\phi$ decay width is determined by the expression

$$\frac{d\Gamma^{\rm BSM}(K^+ \to \pi^+\phi\phi)}{dq^2} = \frac{1}{(2\pi)^3} \cdot \frac{1}{32 M^3_{K_L}} \cdot \frac{2}{M^2_X} \cdot \sqrt{[(M^2_{K^+} + M^2_{\pi^+} - q^2)^2 - 4M^2_{K^+} M^2_{\pi^+}]} \left(1 - \frac{4m^2_\phi}{q^2}\right) \cdot \times \left[\frac{f_0(q^2)(M^2_{K^+} - m^2_{\pi^+})}{2(-m_d + m_s)} \right]^2. \quad (27)$$

It is convenient to use the ratio $\beta^{-1} \equiv \frac{\Gamma(K_L \to \phi\phi)}{\Gamma(K^+ \to \pi^+\nu\bar{\nu})}$, because the ratio β does not depend on the unknown value of M_X or on the values of quark masses m_d, m_s. For the case $m_\phi \ll M_{\pi^+}$ we find that

$$\beta \approx 10^{-2}. \quad (28)$$

As in the previous model, the smallness of the β is mainly due to the three-body phase space smallness in comparison with two-body phase space.

From (16) and (28), we find

$$\mathrm{Br}(K_L \to \phi\phi) \lesssim 2 \times 10^{-8}. \quad (29)$$

For a not very light ϕ particle, the limit on $\mathrm{Br}(K_L \to \phi\phi)$ will be not so stringent as the bound (29); moreover, for ϕ particle mass $M_{K_L}/2 \geq m_\phi \geq (M_{K^+} - M_{\pi^+})/2$, the decay $K^+ \to \pi^+\nu\bar{\nu}$ is kinematically prohibited, while the decay $K_L \to \phi\phi$ is allowed. Therefore, the bound (29) derived from the decay width of the $K^+ \to \pi^+\nu\bar{\nu}$ decay does not work for $K_L \to \phi\phi$ decay mode. Note that such a sub-Gev scalar ϕ could be a good dark matter candidate [23]. As in the previous model, the bound from the $K_L - K_S$ mass difference leads to the bound on the unknown parameter $\delta m^2_{HH_2}$ at the level $\delta m^2_{HH_2} \leq 30\ {\rm GeV}^2$ for $\frac{M_{H_2}}{h_{2d_L s_R}} = 10^4$ GeV.

IV. $K_L \to$ *invisible* DECAY IN MODEL WITH MIRROR WORLD

Finally, we discuss the K_L oscillations into a hidden sector, which would manifest themselves through the $K_L \to invisible$ decay. As an example of such a hidden sector, we consider the one of *the mirror matter models*. The idea that along with ordinary matter may exist its exact mirror copy, introduced for parity conservation, is not new [24]. Accordingly, each ordinary particle of the SM has a corresponding mirror partner of exactly the same mass as the ordinary one. The mirror fields are all singlets under the SM $SU_c(3) \otimes SU_L(2) \otimes U(1)$ gauge group. Mirror matter is dark in terms of the SM interactions, and could be a good candidate for dark matter, see, e.g., Ref. [25], and recently Ref. [26]. In addition to gravity, the interaction between this type of dark matter our own could be transmitted by some gauge singlet particles interacting with both sectors. Any neutral, elementary or composite particle, in principle, can have mixing with its mirror duplicate. This results in several interesting phenomena, such as Higgs [27], positronium [28], muonium [29], or neutron [30] oscillations into their hidden partner, which have been or are planned to be experimentally tested [31–34].

In particular, the neutral K_L meson can mix with it mirror (m) analog $K_{L,m}$ due to effective four-fermion interaction

$$L_{\text{int}} = \frac{1}{M_m^2}[\bar{d}\gamma^\mu(1-\gamma_5)s\bar{s}_m\gamma_\mu(1-\gamma_5)d_m]. \tag{30}$$

The interaction (30) leads to conversion of ordinary K_L mesons to mirror K_L mesons. The decays of mirror K_L mesons are invisible in our world, leading to invisible K_L decay with the branching ratio

$$\text{Br}(K_L \to invisible) = \frac{\delta^2}{2(\delta^2 + \Gamma_{\text{tot}}^2(K_L))}, \tag{31}$$

where

$$\delta = \frac{1}{M_{K_L}}\langle K_{L,m}|L_{\text{int}}|K_L\rangle. \tag{32}$$

For the interaction (27) in the vacuum insertion approximation, we find that

$$\delta \approx \frac{F_K^2 M_{K_L}}{M_m^2}. \tag{33}$$

Numerically, for $\text{Br}(K_L \to invisible) = 10^{-6}$, we can probe the value of M_m up to

$$M_m \lesssim 8.4 \times 10^8 \text{ GeV}. \tag{34}$$

In our estimates we used nonrenormalizable effective four-fermion interaction (30). It is possible to obtain the effective interaction (30) from a renormalizable mirror world model with the Higgs doublet extension of the SM model (see previous discussions) and with the additional interaction term between our and mirror world

$$L_m = \lambda_m(H^+H_2)(H_m H_{m,2}^+) + \text{H.c.} \tag{35}$$

After electroweak symmetry breaking in both our own and the mirror world ($\langle H\rangle = \langle H_m\rangle \approx 174$ GeV), we find an effective four-fermion interaction

$$L_{\text{eff}} = \frac{1}{M_m^2}\bar{d}_L s_R \bar{s}_{R,m} d_{L,m} + \text{H.c.}, \tag{36}$$

where

$$\frac{1}{M_m^2} = \frac{h_{2d_L s_R}^2}{M_{H_2}^2} \cdot \frac{\lambda_m|\langle H\rangle|^2}{M_{H_2}^2}. \tag{37}$$

V. CONCLUSION

In conclusion, the observation of the $K_L \to invisible$ decay with the branching ratio $\text{Br}(K_L \to invisible) \gg 10^{-10}$ would unambiguously signal the presence of BSM physics. We consider the $K_L \to invisible$ decay in several natural extensions of the SM, such as the 2HDM, 2HDM and light neutral scalar field ϕ, and mirror dark matter model. Using constraints from the experimental value for $\text{Br}(K^+ \to \pi^+\nu\bar{\nu})$, we find that the $K_L \to invisible$ decay branching ratio could be in the region $\text{Br}(K_L \to invisible) \simeq 10^{-8}$–$10^{-6}$, which is experimentally accessible, allowing us to test new physics scales well above 100 TeV. In some scenarios these bounds can be avoided, as in the model with the massive right-handed neutrino and scalar ϕ particle. This makes the $K_L \to invisible$ decay a powerful clean probe of new physics, that is complementary to other rare K decay channels. Additionally, in the case of observation, the $K_L \to invisible$ decay could influence the precision of the Bell-Steinberger analysis of the $K^0 - \bar{K}^0$ system. The results obtained provide a strong motivation for a sensitive search for this process in a near-future K decay experiment proposed in Ref. [9]. It should be noted that in full analogy with the case of K_L invisible decay, we can expect the existence of invisible decays of B_d and B_s mesons, see e.g. Refs. [35,36], with the branchings similar to those discussed above.

ACKNOWLEDGMENTS

This work was supported by Russian Foundation for Basic Research Grant No. 13-02-00363.

[1] A. J. Buras and J. Girrbach, Rep. Prog. Phys. **77**, 086201 (2014).
[2] V. Cirigliano, G. Ecker, H. Neufeld, A. Pich, and J. Portoles, Rev. Mod. Phys. **84**, 399 (2012).
[3] D. Bryman, W. J. Marciano, R. Tschirhart, and T. Yamanaka, Annu. Rev. Nucl. Part. Sci. **61**, 331 (2011).
[4] T. K. Komatsubara, Prog. Part. Nucl. Phys. **67**, 995 (2012).
[5] M. Blanke, *Proc. Sci.*, KAON13 (2013) 010.
[6] L. G. Landsberg, Usp. Fiz. Nauk **173**, 1025 (2003) [Phys. Usp. **46**, 995 (2003)].
[7] K. A. Olive *et al.* (Particle Data Group), Chin. Phys. C **38**, 090001 (2014).
[8] W. J. Marciano and Z. Parsa, Phys. Rev. D **53**, R1 (1996).
[9] S. N. Gninenko, Phys. Rev. D **91**, 015004 (2015).
[10] M. Antonelli and G. D′Ambrosio, *CPT* invariance tests in neutral kaon decay, p. 920, in Ref. [7].
[11] K. R. Schubert, L. Li Gioi, A. J. Bevan, and A. Di Domenico, arXiv:1401.6938.
[12] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher, and J. P. Silva, Phys. Rep. **516**, 1 (2012).
[13] In Ref. [14], a model with an additional Higgs isodoublet interacting only with lepton fields was proposed. In this model the neutrino acquires nonzero Dirac masses due to the nonzero vacuum expectation value of the second Higgs isodoublet that allows the seesaw scale to decrease from $O(10^{15})$ GeV to $O(10^3)$ GeV.
[14] E. Ma, Phys. Rev. Lett. **86**, 2502 (2001).
[15] N. V. Krasnikov and V. F. Tokarev, Yad. Fiz. **33**, 283 (1981).
[16] The quark masses m_d, m_s and the effective mass M_X implicitly depend on the renormalization point μ, but their combination $M_X^2(m_d + m_s)$ and hence the decay width (7) is renormalization-group invariant and does not depend on the renormalization point μ.
[17] In our estimate (8), we used the values $\tau(K_L) = 5.17 \times 10^{-8}$ sec, $m_{\nu_\tau} = 0$ and $(m_d + m_s)(\mu = 1\ \mathrm{GeV}) = 160$ MeV.
[18] V. V. Artamonov *et al.*, Phys. Rev. Lett. **101**, 191802 (2008).
[19] J. Ahn *et al.*, Phys. Rev. D **81**, 072004 (2010).
[20] P. A. Boyle, J. M. Flynn, A. Jüttner, C. Kelly, C. Maynard, H. Pedroso de Lima, C. T. Sachrajda, and J. M. Zanotti, Eur. Phys. J. C **69**, 159 (2010).
[21] D. Cutts, R. Stiening, C. Wiegand, and M. Deutsch, Phys. Rev. **184**, 1380 (1969).
[22] As a review, see for example J. N. Butler *et al.*, arXiv:1311.1076.
[23] C. Boehm and P. Fayet, Nucl. Phys. **B683**, 219 (2004).
[24] T. D. Lee and C. N. Yang, Phys. Rev. **104**, 254 (1956); I. Kobzarev, L. Okun, and I. Pomeranchuk, Sov. J. Nucl. Phys. **3**, 837 (1966).
[25] S. I. Blinnikov and M. Yu. Khlopov, Sov. J. Nucl. Phys. **36**, 472 (1982); Sov. Astron. J. **27**, 371 (1982); R. Foot, H. Lew, and R. R. Volkas, Phys. Lett. B **272**, 67 (1991); Z. Berezhiani and R. Mohapatra, Phys. Rev. D **52**, 6607 (1995); L. B. Okun, Usp. Fiz. Nauk **177**, 397 (2007) [Phys. Usp. **50**, 380 (2007)].
[26] R. Foot, Int. J. Mod. Phys. A **29**, 1430013 (2014).
[27] A. Yu. Ignatiev and R. R. Volkas, Phys. Lett. B **487**, 294 (2000); B. Patt and F. Wilczek, arXiv:hep-ph/0605188.
[28] S. L. Glashow, Phys. Lett. **167B**, 35 (1986); S. N. Gninenko, Phys. Lett. B **326**, 317 (1994); R. Foot and S. N. Gninenko, Phys. Lett. B **480**, 171 (2000).
[29] S. N. Gninenko, N. V. Krasnikov, and V. A. Matveev, Phys. Rev. D **87**, 015016 (2013).
[30] Z. Berezhiani and L. Bento, Phys. Rev. Lett. **96**, 081801 (2006); Z. Berezhiani, Eur. Phys. J. C **64**, 421 (2009).
[31] A. Badertscher, P. Crivelli, U. Gendotti, S. N. Gninenko, V. Postoev, A. Rubbia, V. Samoylenko, and D. Sillou, Phys. Rev. D **75**, 032004 (2007).
[32] G. Banet *et al.*, Phys. Rev. Lett. **99**, 161603 (2007).
[33] A. P. Serebrov *et al.*, Phys. Lett. B **663**, 181 (2008).
[34] P. Crivelli, A. Belov, U. Gendotti, S. Gninenko, and A. Rubbia, JINST **5**, P08001 (2010).
[35] A. Badin and A. Petrov, Phys. Rev. D **82**, 034005 (2010).
[36] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. Lett. **93**, 091802 (2004).