Improved constraints on sterile neutrinos in the MeV to GeV mass range

D. A. Bryman[1,2] and R. Shrock[3]

[1]*Department of Physics and Astronomy, University of British Columbia, Vancouver, British Columbia V6T 1Z1, Canada*

[2]*TRIUMF, 4004 Wesbrook Mall, Vancouver, British Columbia V6T 2A3, Canada*

[3]*C. N. Yang Institute for Theoretical Physics and Department of Physics and Astronomy, Stony Brook University, Stony Brook, New York 11794, USA*

(Received 14 April 2019; published 16 September 2019)

Improved upper bounds are presented on the coupling $|U_{e4}|^2$ of an electron to a sterile neutrino ν_4 from analyses of data on nuclear and particle decays, including superallowed nuclear beta decays, the ratios $R^{(\pi)}_{e/\mu} = \mathrm{BR}(\pi^+ \to e^+\nu_e)/\mathrm{BR}(\pi^+ \to \mu^+\nu_\mu)$, $R^{(K)}_{e/\mu}$, $R^{(D_s)}_{e/\tau}$, and B^+_{e2} decay, covering the mass range from MeV to GeV.

DOI: 10.1103/PhysRevD.100.053006

Neutrino oscillations and hence neutrino masses and lepton mixing have been established and are of great importance as physics beyond the original Standard Model (SM). Most oscillation experiments with solar, atmospheric, accelerator, and reactor (anti)neutrinos [1–10] can be explained within the minimal framework of three neutrino mass eigenstates with values of $\Delta m^2_{ij} = m^2_{\nu_i} - m^2_{\nu_j}$ given approximately by $\Delta m^2_{21} = 0.74 \times 10^{-4}$ eV^2 and $|\Delta m^2_{32}| = 2.5 \times 10^{-3}$ eV^2, with normal mass ordering $m_{\nu_3} > m_{\nu_2}$ favored; furthermore, the lepton mixing angles θ_{23}, θ_{12}, and θ_{13} have been measured, with a tentative indication of a nonzero value of the CP-violating quantity $\sin(\delta_{CP})$ (for compilations and fits, see [11–16]).

In addition to the three known neutrino mass eigenstates, there could be others, which would necessarily be primarily electroweak singlets (sterile) [17] (see, e.g., [18]). Indeed, sterile neutrinos are present in many ultraviolet (UV) extensions of the SM. Whether sterile neutrinos exist in nature is one of the most outstanding questions in particle physics, and therefore, improved constraints on their couplings are of fundamental and far-reaching importance. Taking account of the possibility of sterile neutrinos, the neutrino interaction eigenstates ν_ℓ would be given by

$$\nu_\ell = \sum_{i=1}^{3+n_s} U_{\ell i}\nu_i, \tag{1}$$

where $\ell = e, \mu, \tau$; n_s denotes the number of sterile neutrinos; and U is the lepton mixing matrix [19].

Here we obtain improved upper limits on $|U_{ei}|^2$ for a sterile neutrino ν_i in a wide range of masses from the MeV to GeV scale and point out new experiments that would be worthwhile and could yield further improvements. For simplicity, we assume one heavy neutrino, $n_s = 1$, with $i = 4$; it is straightforward to generalize to $n_s \geq 2$. Since a ν_4 in this mass range decays, it is not excluded by the cosmological upper limit on the sum of effectively stable neutrinos, $\sum_i m_{\nu_i} \lesssim 0.12$ eV [20]. Such a ν_4 is subject to a number of constraints from cosmology (e.g., [21]); however, since these depend on assumptions about the early universe, we choose here to focus on direct laboratory bounds. Constraints from the nonobservation of neutrinoless double beta decay are satisfied by assuming that ν_4 is a Dirac neutrino [22]. Since sterile neutrinos violate the conditions for the diagonality of the weak neutral current [23,24], ν_4 has invisible tree-level decays of the form $\nu_4 \to \nu_j\bar{\nu}_i\nu_i$ where $1 \leq i,\ j \leq 3$ with model-dependent invisible branching ratios. Because our bounds are purely kinematic, they are complementary to bounds from searches for neutrino decays, which involve model-dependent assumptions on branching ratios into visible versus invisible final states.

We first obtain improved upper bounds on $|U_{e4}|^2$ from nuclear beta decay data. The emission of a ν_4 via lepton mixing in nuclear beta decay has several effects, including producing a kink in the Kurie plot and reducing the decay rate [25]. For the nuclear beta decays $(Z, A) \to (Z + 1, A) + e^- + \bar{\nu}_e$ or $(Z, A) \to (Z - 1, A) + e^+ + \nu_e$ into a set of neutrino mass eigenstates $\nu_i \in \nu_e$, $i = 1, 2, 3$ of negligibly small masses, plus a ν_4 of non-negligible mass, the differential decay rate is

$$\frac{dN}{dE} = C[(1-|U_{e4}|^2)pE(E_0-E)^2 + |U_{e4}|^2 pE(E_0-E)[(E_0-E)^2 - m_{\nu_4}^2]^{1/2} \times \theta(E_0 - E - m_{\nu_4})], \tag{2}$$

where $p \equiv |\mathbf{p}|$ and E denote the 3-momentum and (total) energy of the outgoing $e^{\pm}$, E_0 denotes its maximum energy for the SM case, the Heaviside θ function is defined as $\theta(x) = 1$ for $x > 0$ and $\theta(x) = 0$ for $x \le 0$, and $C = G_F^2 |V_{ud}|^2 F_F |\mathcal{M}|^2/(2\pi^3)$, where $\mathcal{M}$ denotes the nuclear transition matrix element, V is the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix, and F_F is the Fermi function. Early bounds on $|U_{e4}|^2$ were set from searches for kinks in Kurie plots in [25] and analyses of particle decays [26–28], and from dedicated experiments. For example, a search for kinks in the Kurie plot in ^{20}F beta decay reported in Ref. [29] yielded an upper bound on $|U_{e4}|^2$ decreasing from 5.9×10^{-3} for $m_{\nu_4} = 0.4$ MeV to 1.8×10^{-3} for $m_{\nu_4} = 2.8$ MeV. (Some recent reviews of searches for sterile neutrinos include [30–35].)

In addition to kink searches, a powerful method to set constraints on massive neutrino emission, via lepton mixing, in nuclear beta decays is to analyze the decay rates. Since, in general, the heavy neutrino would also be emitted in μ decay, the measurement of the μ lifetime performed assuming the SM would yield an apparent (app) value of the Fermi constant, denoted $G_{F,\mathrm{app}}$, that would be smaller than the true value, G_F, given at tree level by $G_F/\sqrt{2} = g^2/(8m_W^2)$, where g is the SU(2) gauge coupling [26–28]. To avoid this complication, the ratios of rates of different nuclear beta decays are compared.

The integration of dN/dE over E gives the kinematic rate factor f. The combination of this with the half-life for the nuclear beta decay, $t \equiv t_{1/2}$, yields the product ft. Incorporation of nuclear and radiative corrections yields the corrected ft value for a given decay, denoted $\mathcal{F}t$. Conventionally, analyses of the $\mathcal{F}t$ values for the most precisely measured superallowed $0^+ \to 0^+$ nuclear beta decays have been used, in conjunction with the value of $G_{F,\mathrm{app}}$ from μ decay, to infer a value of the weak mixing matrix element, $|V_{ud}|$ [36–45]. A first step in these analyses has been to establish the mutual consistency of the $\mathcal{F}t$ values for these superallowed $0^+ \to 0^+$ decays. Since the emission of a ν_4 with mass of a few MeV would have a different effect on the kinematic functions and integrated rates for nuclear beta decays with different Q (energy release) values, it would upset this mutual consistency.

Hence, from this mutual agreement of $\mathcal{F}t$ values, an upper limit on $|U_{e4}|^2$ can be derived for values of m_{ν_4} such that a ν_4 could be emitted in some of these superallowed decays. Reference [37] obtained upper bounds on $|U_{e4}|^2$ ranging from 10^{-2} down to 2×10^{-3} for m_{ν_4} from 0.5 to 2 MeV, while Ref. [29] obtained the

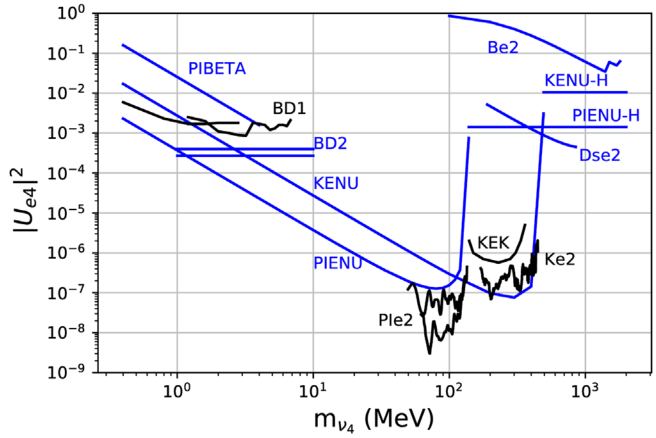

FIG. 1. The 90% C.L. upper limits on $|U_{e4}|^2$ vs m_{ν_4} from various sources: PIBETA, pion beta decay BD1, previous limits from beta decay [29]; BD2, beta decay with the two dashed horizonal lines based on our analysis using [42] and [43]; PIENU and PIENU-H, the ratio $\frac{\mathrm{BR}(\pi^+\to e^+\nu_e)}{\mathrm{BR}(\pi^+\to \mu^+\nu_\mu)}$ in the kinematically allowed and forbidden regions for ν_4 emission; $Pie2$, $\pi^+ \to e^+\nu_{e4}$ peak search [47]; KENU and KENU-H, the ratio $\frac{\mathrm{BR}(K^+\to e^+\nu_e)}{\mathrm{BR}(K^+\to \mu^+\nu_\mu)}$ in the kinematically allowed and forbidden regions for ν_4 emission; $Ke2$, $K^+ \to e^+\nu_{e4}$ peak search [48]; $Dse2$, $D_s^+ \to e^+\nu_{e4}$; and $Be2$, $B^+ \to e^+\nu_{e4}$.

limits $|U_{e4}|^2 < 1 \times 10^{-3}$ to $|U_{e4}|^2 < 2 \times 10^{-3}$ depending nonmonotonically on m_{ν_4} from 1 to 7 MeV. The maximum Q value in the current set of 14 superallowed $0^+ \to 0^+$ beta decays used for the $\mathcal{F}t$ fit in [41,42] is 9.4 MeV (for ^{74}Rb). A measure of the mutual agreement is the precision with which $|V_{ud}|^2$ is determined, so a reduction in the fractional uncertainty of the value of $|V_{ud}|^2$ results in an improved upper limit on $|U_{e4}|^2$. Reference [37] obtained $|V_{ud}| = 0.9740 \pm 0.001$. The recent analyses in [42] and [43] obtained $|V_{ud}| = 0.97420(21)$ and $|V_{ud}| = 0.97370(14)$, respectively [46]. Applying these factors of improvement from [42] and [43] to the previous bounds in [37], improved upper bounds are obtained as

$$|U_{e4}|^2 < 4 \times 10^{-4} \tag{3}$$

and

$$|U_{e4}|^2 < 2.7 \times 10^{-4} \tag{4}$$

for ν_4 masses in the range $1\ \mathrm{MeV} \lesssim m_{\nu_4} < 9.4$ MeV, indicated in Fig. 1 as BD2. (These and other limits presented are at the 90% confidence level.)

We next discuss upper bounds from two-body leptonic decays of charged pseudoscalar mesons (generically denoted as M^+) [25,26]. This method is quite powerful because the signal is a monochromatic peak in dN/dp_ℓ, and for $M^+ \to e^+\nu_e$ (denoted M_{e2}^+) decays, the strong helicity suppression in the SM case is removed when a heavy neutrino is emitted. The presence of a massive ν_4 also changes the ratio of branching ratios $\mathrm{BR}(M^+ \to e^+\nu_e)/\mathrm{BR}(M^+ \to \mu^+\nu_\mu)$ from its SM value, and this was used to set further bounds [25,26,49]. A number of dedicated experiments have been performed to search for a peak due to heavy neutrino emission and also to measure $\mathrm{BR}(M^+ \to e^+\nu_e)/\mathrm{BR}(M^+ \to \mu^+\nu_\mu)$ with $\pi^+_{\ell 2}$, $K^+_{\ell 2}$, and $B^+_{\ell 2}$, where $\ell = e, \mu$ [47–60].

In the SM with only the three known neutrinos with negligibly small masses, the ratio

$$R^{(M)}_{\ell/\ell'} \equiv \frac{\mathrm{BR}(M^+ \to \ell^+\nu_\ell)}{\mathrm{BR}(M^+ \to \ell'^+\nu_{\ell'})} \tag{5}$$

is given by

$$R^{(M)}_{\ell/\ell',\mathrm{SM}} = \frac{m_\ell^2}{m_{\ell'}^2}\left[\frac{1-\delta^{(M)}_\ell}{1-\delta^{(M)}_{\ell'}}\right]^2 (1+\delta_{\mathrm{RC}}), \tag{6}$$

where $\delta^{(M)}_\ell = m_\ell^2/m_M^2$ and δ_{RC} is the radiative correction (RC) [61–66].

We denote the ratio of the experimental measurement of $R^{(M)}_{\ell/\ell'}$ to the SM prediction as

$$\bar{R}^{(M)}_{\ell/\ell'} \equiv \frac{R^{(M)}_{\ell/\ell'}}{R^{(M)}_{\ell/\ell',\mathrm{SM}}}. \tag{7}$$

The most precise measurement of $R^{(\pi)}_{e/\mu}$ is from the PIENU experiment at TRIUMF, with the result $R^{(\pi)}_{e/\mu} = (1.2344 \pm 0.0023_{\mathrm{stat}} \pm 0.0019_{\mathrm{syst}}) \times 10^{-4}$ [58]. Including [67–69], the resultant PDG world average is $R^{(\pi)}_{e/\mu} = (1.2327 \pm 0.0023) \times 10^{-4}$ [11], in agreement with the SM prediction with RC, $R^{(\pi)}_{e/\mu} = (1.2352 \pm 0.0002) \times 10^{-4}$ [62,63,65], resulting in

$$\bar{R}^{(\pi)}_{e/\mu} = 0.9980 \pm 0.0019. \tag{8}$$

The ratio $R^{(K)}_{e/\mu}$ has recently been measured by the NA62 experiment at CERN [56], dominating the world average [11]

$$R^{(K)}_{e/\mu} = (2.488 \pm 0.009) \times 10^{-5}. \tag{9}$$

The SM prediction with RC [63,66] is

$$R^{(K)}_{e/\mu,\mathrm{SM}} = (2.477 \pm 0.001) \times 10^{-5}, \tag{10}$$

resulting in

$$\bar{R}^{(K)}_{e/\mu} = 1.0044 \pm 0.0037. \tag{11}$$

With emission of a heavy neutrino ν_4, the ratio $R^{(M)}_{\ell/\ell',\mathrm{SM}}$ for general ℓ, ℓ' changes to

$$R^{(M)}_{\ell/\ell'} = \left[\frac{[(1-|U_{\ell 4}|^2)\rho(\delta^{(M)}_\ell, 0) + |U_{\ell 4}|^2\rho(\delta^{(M)}_\ell, \delta^{(M)}_{\nu_4})]}{(1-|U_{\ell' 4}|^2)\rho(\delta^{(M)}_{\ell'}, 0) + |U_{\ell' 4}|^2\rho(\delta^{(M)}_{\ell'}, \delta^{(M)}_{\nu_4})}\right] \times (1+\delta_{\mathrm{RC}}), \tag{12}$$

where $\delta^{(M)}_{\nu_4} = m_{\nu_4}^2/m_M^2$, and the kinematic function $\rho(x,y)$ is [25,26]

$$\rho(x,y) = [x + y - (x-y)^2][\lambda(1,x,y)]^{1/2}, \tag{13}$$

with

$$\lambda(z,x,y) = x^2 + y^2 + z^2 - 2(xy + yz + zx). \tag{14}$$

Thus, in the SM case, $\rho(x,0) = x(1-x)^2$. Here and below, it is implicitly understood that $\rho(\delta^{(M)}_\ell, \delta^{(M)}_{\nu_4}) = 0$ if $m_{\nu_4} \geq m_M - m_\ell$, where the decay $M^+ \to \ell^+\nu_4$ is kinematically forbidden. We define

$$\bar{\rho}(x,y) = \frac{\rho(x,y)}{\rho(x,0)} = \frac{\rho(x,y)}{x(1-x)^2} \tag{15}$$

so

$$\bar{R}^{(M)}_{\ell/\ell'} = \frac{1 - |U_{\ell 4}|^2 + |U_{\ell 4}|^2\bar{\rho}(\delta^{(M)}_\ell, \delta^{(M)}_{\nu_4})}{1 - |U_{\ell' 4}|^2 + |U_{\ell' 4}|^2\bar{\rho}(\delta^{(M)}_{\ell'}, \delta^{(M)}_{\nu_4})}. \tag{16}$$

With no loss of generality, we order ℓ and ℓ' such that $m_{\ell'} > m_\ell$ and define the mass intervals (i) $I^{(M)}_1 \colon m_{\nu_4} < m_M - m_{\ell'}$; (ii) $I^{(M)}_2 \colon m_M - m_{\ell'} < m_{\nu_4} < m_M - m_\ell$; and (iii) $I^{(M)}_3 \colon m_{\nu_4} > m_M - m_\ell$. Thus, a ν_4 with $m_{\nu_4} \in I^{(M)}_1$ contributes to both $M^+_{\ell 2}$ and $M^+_{\ell' 2}$ decays, while if $m_{\nu_4} \in I^{(M)}_2$, then ν_4 contributes to $M^+_{\ell 2}$, but not to $M^+_{\ell' 2}$ decay, and if $m_{\nu_4} \in I^{(M)}_3$, then ν_4 cannot be emitted in either $M^+_{\ell 2}$ or $M^+_{\ell' 2}$ decay.

If for a given m_{ν_4} one knows, e.g., from peak-search experiments, that $|U_{\ell' 4}|^2$ is sufficiently small that the denominator of (16) can be approximated well by 1, then an upper bound on the deviation of $\bar{R}^{(M)}_{\ell/\ell'}$ from 1 yields an upper bound on $|U_{\ell 4}|^2$. Thus, one has the bound

$$|U_{\ell 4}|^2 < \frac{\bar{R}^{(M)}_{\ell/\ell'} - 1}{\bar{\rho}(\delta^{(M)}_\ell, \delta^{(M)}_{\nu_4}) - 1} \quad \text{for } m_{\nu_4} \in I^{(M)}_2. \tag{17}$$

This gives very stringent upper limits on $|U_{\ell 4}|^2$ because $\bar{\rho}(\delta^{(M}_e, \delta^{(M)}_{\nu_4}) \gg 1$ over much of the interval $I^{(M)}_2$

(see Figs. 3–5 in [26]). If $m_{\nu_4} \in I_3^{(M)}$, then (16) reduces to $\bar{R}^{(M)}_{\ell/\ell'} = (1-|U_{\ell 4}|^2)/(1-|U_{\ell' 4}|^2)$, so if $|U_{\ell' 4}|^2 \ll 1$ in this interval, then the upper limit is

$$|U_{\ell 4}|^2 < 1 - \bar{R}^{(M)}_{\ell/\ell'} \quad \text{for } m_{\nu_4} \in I_3^{(M)}. \tag{18}$$

We now apply this analysis to $R^{(\pi)}_{e/\mu}$, using (17) and (18) with $M^+ = \pi^+$, $\ell = e$, and $\ell' = \mu$. From previous $\pi^+_{\mu 2}$ peak search experiments [47,48,50–60] and the calculation of $\bar{\rho}(\delta^{(\pi)}_\mu, \delta^{(\pi)}_{\nu_4})$, it follows that $|U_{\mu 4}|^2$ is sufficiently small for $m_{\nu_4} \in I_2^{(\pi)}$ that we can approximate the denominator of Eq. (16) by 1. From $\bar{R}^{(\pi)}_{e/\mu}$ in Eq. (8), using the procedure from [70], we obtain the limit $\bar{R}^{(\pi)}_{e/\mu} < 1.0014$. Then, for $\nu_4 \in I_2^{(\pi)}$, we find

$$|U_{\ell 4}|^2 < \frac{\bar{R}^{(\pi)}_{e/\mu'} - 1}{\bar{\rho}(\delta^{(\pi)}_e, \delta^{(\pi)}_{\nu_4}) - 1} < \frac{0.0014}{\bar{\rho}(\delta^{(\pi)}_e, \delta^{(\pi)}_{\nu_4}) - 1}. \tag{19}$$

This bound is labeled as PIENU in Fig. 1. If $m_{\nu_4} \in I_3^{(\pi)}$, i.e., $m_{\nu_4} > 139$ MeV, then, using (18), we obtain the upper bound on $|U_{e4}|^2$ given by the flat line labeled PIENU-H in Fig. 1.

We next obtain a bound on $|U_{e4}|^2$ by applying the same type of analysis to $R^{(K)}_{e/\mu}$. From $K_{\mu 2}$ peak search experiments [48,51,57] and the calculation of $\bar{\rho}(\delta^{(K)}_\mu, \delta^{(K)}_{\nu_4})$, $|U_{\mu 4}|^2$ is sufficiently small that we can approximate the denominator of Eq. (16) well by 1. Using Eq. (11) for $\nu_4 \in I_2^{(K)}$, we find

$$|U_{\ell 4}|^2 < \frac{\bar{R}^{(K)}_{e/\mu'} - 1}{\bar{\rho}(\delta^{(K)}_e, \delta^{(K)}_{\nu_4}) - 1} < \frac{0.010}{\bar{\rho}(\delta^{(K)}_e, \delta^{(K)}_{\nu_4}) - 1}. \tag{20}$$

This upper limit on $|U_{e4}|^2$ is labeled KENU in Fig. 1. For $m_{\nu_4} \in I_3^{(K)}$, i.e., $m_{\nu_4} > 493$ MeV, using ((18), we obtain the flat upper bound labeled KENU-H in Fig. 1.

One can also apply these methods to two-body leptonic decays of heavy-quark hadrons. We first consider $D_s^+ \to \ell^+ \nu_\ell$ decays [71], using (17) and (18) with $M^+ = D_s^+$, $\ell = e$, and $\ell' = \tau$. Experimental data from CLEO, *BABAR*, Belle, and BES have determined $\mathrm{BR}(D_s^+ \to \mu^+ \nu_\mu) = (5.49 \pm 0.17) \times 10^{-3}$ and $\mathrm{BR}(D_s^+ \to \tau^+ \nu_\tau) = (5.48 \pm 0.23) \times 10^{-2}$ [72–76]. Furthermore, searches by CLEO [72], *BABAR* [73], and Belle [74] have yielded the limit $\mathrm{BR}(D_s^+ \to e^+ \nu_e) < 0.83 \times 10^{-4}$. Hence, $R^{(D_s)}_{e/\tau} < 1.6 \times 10^{-3}$. For $R^{(D_s)}_{e/\tau}$, using the results of [62], we calculate $1 + \delta_{\mathrm{RC}} = 0.948$. Substituting this in Eq. (6) with $M = D_s$, $\ell = e$, $\ell' = \tau$, we find

$$R^{(D_s)}_{e/\tau,\mathrm{SM}} = 2.29 \times 10^{-6}. \tag{21}$$

Therefore, $\bar{R}^{(D_s)}_{e/\tau} < 7.0 \times 10^2$. For $R^{(D_s)}_{e/\tau}$, the interval $I_2^{(D_s)}$ is 191 MeV $< m_{\nu_4} <$ 1.457 GeV. Actually, we restrict m_{ν_4} to a lower-mass subset of this interval because for sufficiently great m_{ν_4}, even though the $D_s^+ \to e^+ \nu_4$ decay is kinematically allowed to occur, the momentum p_e (in the D_s rest frame) would be below the minimal value set by experimental cuts in the BES III event reconstruction. With $p_{e,\mathrm{cut}} \simeq 0.8$ GeV [77], this means that m_{ν_4} must be less than 0.85 GeV for the event to be accepted. Thus, we consider 0.191 GeV $< m_{\nu_4} <$ 0.85 GeV. Substituting the experimental limit on $\bar{R}^{(D_s)}_{e/\tau}$ in the special case of (16) with $M = D_s$, $\ell = e$, $\ell' = \tau$ and using the fact that $|U_{\tau 4}|^2 \ll 1$ for this m_{ν_4} mass range [11], we obtain a resultant limit from (17). For $m_{\nu_4} = 0.191$ GeV, $\bar{\rho}(\delta^{(D_s)}_e, \delta^{(D_s)}_{\nu_4}) = 1.37 \times 10^5$, increasing to $\bar{\rho}(\delta^{(D_s)}_e, \delta^{(D_s)}_{\nu_4}) = 1.83 \times 10^6$ for $m_{\nu_4} = 0.85$ GeV. We thus obtain the upper bound on $|U_{e4}|^2$ labeled D_{se2} in Fig. 1.

A dedicated peak-search experiment to search for the heavy-neutrino decay $D_s^+ \to e^+ \nu_4$ would be worthwhile and could improve the upper bound on $|U_{e4}|^2$. Similarly, a search for leptonic D decays like $D^+ \to e^+ \nu_4$ would be valuable and will be discussed elsewhere. The very large values of $\bar{\rho}(\delta^{(D_s)}_e, \delta^{(D_s)}_{\nu_4})$ and $\bar{\rho}(\delta^{(D)}_e, \delta^{(D)}_{\nu_4})$ over a large portion of the kinematically allowed ranges of m_{ν_4} in $D_s^+ \to e^+ \nu_4$ and $D \to e^+ \nu_4$ mean that there would be quite strong kinematic enhancement of the heavy neutrino decay relative to the corresponding $(D_s^+)_{e2}$ and D^+_{e2} decays. In particular, these searches could be performed by the BES III experiment, which recently reported results from a data sample of 3.19 fb^{-1} and expects to collect considerably higher statistics.

Finally, we consider $B^+ \to \ell^+ \nu_\ell$ decays. There is an upper limit $\mathrm{BR}(B^+ \to e^+ \nu_e) < 0.98 \times 10^{-6}$ from Belle [78] and *BABAR* [79]. For the other two leptonic decay modes, $\mathrm{BR}(B^+ \to \mu^+ \nu_\mu) = (6.46 \pm 2.22_{\mathrm{stat}} \pm 1.60_{\mathrm{syst}}) \times 10^{-7}$ from Belle [80], with a recent update $\mathrm{BR}(B^+ \to \mu^+ \nu_\mu) = (5.3 \pm 2.0_{\mathrm{stat}} \pm 0.9_{\mathrm{syst}}) \times 10^{-7}$ [81,82], and $\mathrm{BR}(B^+ \to \tau^+ \nu_\tau) = (1.09 \pm 0.24) \times 10^{-4}$ from *BABAR* [83] and Belle [84,85]. The measured values of $\mathrm{BR}(B^+ \to \mu^+ \nu_\mu)$ are in agreement with the SM prediction $\mathrm{BR}(B^+ \to \mu^+ \nu_\mu)_{\mathrm{SM}} = (3.80 \pm 0.31) \times 10^{-7}$ [80]. The measured value of $\mathrm{BR}(B^+ \to \tau^+ \nu_\tau)$ is also in agreement with the SM prediction $\mathrm{BR}(B^+ \to \tau^+ \nu_\tau)_{\mathrm{SM}} = (0.75^{+0.10}_{-0.05}) \times 10^{-4}$ [85,86].

We focus on data from a $B^+_{\ell 2}$ peak search experiment by Belle [59]. In general [25],

$$\frac{\mathrm{BR}(M^+ \to \ell^+ \nu_4)}{\mathrm{BR}(M^+ \to \ell^+ \nu_\ell)_{\mathrm{SM}}} = \frac{|U_{\ell 4}|^2 \bar{\rho}(\delta^{(M)}_\ell, \delta^{(M)}_{\nu_4})}{1 - |U_{\ell 4}|^2}. \tag{22}$$

For m_{ν_4} in the range from 0.1 GeV to 1.4 GeV, the Belle experiment obtained an upper limit on $\mathrm{BR}(B^+ \to e^+\nu_4)$ of 2.5×10^{-6}, while in the interval of m_{ν_4} from 1.4 GeV to 1.8 GeV, this upper limit increased to 7×10^{-6}. In the range of m_{ν_4} from 0.1 to 1.3 GeV, the Belle experiment obtained (nonmonotonic) upper limits on $\mathrm{BR}(B^+ \to \mu^+\nu_4)$ of approximately 2–4 $\times 10^{-6}$, and in the interval of m_{ν_4} from 1.3 GeV to 1.8 GeV, it obtained upper limits varying from 2×10^{-6} to 1.1×10^{-5}. Substituting the $\mathrm{BR}(B^+ \to e^+\nu_4)$ limits in Eq. (22) with $M = B$ and $\ell = e$, we obtain the upper limits on $|U_{e4}|^2$ shown as the curve B_{e2} in Fig. 1 [87]. From the $\mathrm{BR}(B^+ \to e^+\nu_4)$ limits we infer upper limits on $|U_{\mu4}|^2$ that decrease from 0.83 to 3.4×10^{-2} as m_{ν_4} increases from 0.1 GeV to 1.2 GeV. Further peak searches for $B^+ \to \ell^+\nu_4$ with $\ell = e,\ \mu$ at Belle II would be worthwhile as a higher-statistics extension of [59].

We briefly remark on other constraints on a Dirac ν_4 in the mass range considered here. From the results of [23,88], it follows that there is a negligibly small contribution to decays such as $\mu \to e\gamma$ and $\mu \to ee\bar{e}$. Similarly, there is no conflict with bounds on neutrino magnetic moments [11,89], and contributions to invisible Higgs decays [90] are well below the current upper limit of $\mathrm{BR}(H \to \mathrm{invis}) < 19\%$ [91].

In this work, improved upper limits on $|U_{e4}|^2$ have been presented covering most of the range from $m_{\nu_4} \simeq 1$ MeV to $m_{\nu_4} \simeq 1$ GeV, representing the best available laboratory bounds for a Dirac neutrino ν_4 that do not make model-dependent assumptions concerning visible neutrino decay modes. Over parts of this range, the bounds obtained are competitive with those that assume specific visible ν_4 decays. For example, for $m_{\nu_4} = 30$ MeV, our upper bound is $|U_{e4}|^2 < 0.8 \times 10^{-6}$, while the best bound for this value of m_{ν_4} from experiments searching for neutrino decays is $|U_{e4}|^2 < 1 \times 10^{-6}$ [92]. New peak search experiments to search for $D_s^+ \to e^+\nu_4$ and $D^+ \to e^+\nu_4$ as well as a continued search for $B^+ \to e^+\nu_4$ and continued searches for $\pi^+ \to e^+\nu_4$ [47] and $K^+ \to e^+\nu_4$ [48] would be valuable; these could improve the bounds further. Other constraints on sterile neutrinos, such as from $\pi^+ \to \pi^0 e^+\nu_e$ decay, and a detailed report of the results presented here will be published elsewhere.

ACKNOWLEDGEMENTS

We thank J. Benitez, J. Hardy, V. Luth, W. Marciano, M. Ramsey-Musolf, C. Yuan, and G. Zhao for useful discussions. This work was supported in part by the Natural Sciences and Engineering Research Council and TRIUMF through a contribution from the National Research Council of Canada (D. B.) and by the U.S. National Science Foundation Grant No. NSF-PHY-16-1620628 (R. S.).

[1] R. Davis, D. S. Harmer, and K. C. Hoffman, Phys. Rev. Lett. **20**, 1205 (1968).
[2] A. I. Abazov (SAGE Collaboration), Phys. Rev. Lett. **67**, 3332 (1991); J. N. Abdurashitov *et al.* (SAGE Collaboration), Phys. Lett. B **328**, 234 (1994); P. Anselmann *et al.* (GALLEX Collaboration), Phys. Lett. B **285**, 376 (1992); Phys. Lett. B **327**, 377 (1994).
[3] Y. Fukuda *et al.* (Kamiokande Collaboration), Phys. Rev. Lett. **77**, 1683 (1996).
[4] Y. Fukuda *et al.* (Super-Kamiokande Collaboration), Phys. Rev. Lett. **81**, 1562 (1998).
[5] Q. R. Ahmad *et al.* (SNO Collaboration), Phys. Rev. Lett. **89**, 011301 (2002).
[6] K. Eguchi *et al.* (KamLAND Collaboration), Phys. Rev. Lett. **90**, 021802 (2003).
[7] M. H. Ahn *et al.* (K2K Collaboration), Phys. Rev. Lett. **90**, 041801 (2003).
[8] C. Arpesella *et al.* (Borexino Collaboration), Phys. Rev. Lett. **101**, 091302 (2008).
[9] Y. Abe *et al.* (Double Chooz Collaboration), Phys. Rev. Lett. **108**, 131801 (2012); F. P. An *et al.* (Daya Bay Collaboration), Phys. Rev. Lett. **108**, 171803 (2012); J. K. Ahn *et al.* (RENO Collaboration), Phys. Rev. Lett. **108**, 191802 (2012).
[10] P. Adamson *et al.* (MINOS Collaboration), Phys. Rev. Lett. **107**, 181802 (2011); K. Abe *et al.* (T2K Collaboration), Phys. Rev. Lett. **112**, 061802 (2014); P. Adamson *et al.* (NOvA Collaboration), Phys. Rev. Lett. **116**, 151806 (2016).
[11] M. Tanabashi *et al.* (Particle Data Group and PDG Collaboration), Phys. Rev. D **98**, 030001 (2018), online at http://pdg.lbl.gov.
[12] See http://www.nu-fit.org.
[13] See proceedings at Neutrino 2018, Heidelberg, https://www.mpi-hd.mpg.de/nu2018.
[14] S. Gariazzo, C. Giunti, M. Laveder, and Y. F. Li, Phys. Lett. B **782**, 13 (2018).
[15] M. Dentler, A. Hernández-Cabezudo, J. Kopp, P. Machado, I. Martinez-Soler, and T. Schwetz, J. High Energy Phys. 08 (2018) 010.
[16] I. Esteban, M. C. Gonzalez-Garcia, A. Hernández-Cabezudo, M. Maltoni, and T. Schwetz, J. High Energy Phys. 01 (2019) 106.
[17] We will use the term “sterile neutrino” both in its precise sense as an electroweak-singlet interaction eigenstate and in a commonly used approximate sense as the corresponding mainly sterile mass eigenstate(s) in this neutrino interaction eigenstate.

[18] A. A. Aguilar-Arevalo *et al.* (MiniBooNE Collaboration), Phys. Rev. Lett. **105**, 181801 (2010); **121**, 221801 (2018).

[19] See, e.g., [14,15] for global fits to neutrino data that allow for a possible sterile neutrino with mass $\sim O(1)$ eV^2.

[20] N. Aghanim *et al.* (Planck Collaboration), arXiv: 1807.06209; P. A. R. Ade *et al.* (Planck Collaboration), Astron. Astrophys. **594**, A13 (2016).

[21] There is extensive literature on these constraints and on models that are designed to evade one or more of them. See, e.g., the following and references therein: S. Hannestad and G. G. Raffelt, Phys. Rev. D **72**, 103514 (2005); G. Gelmini, E. Osoba, S. Palomares-Ruiz, and S. Pascoli, J. Cosmol. Astropart. Phys. 10 (2008) 029; G. Steigman, Adv. HEP **2012**, 268321 (2012); J. Lesgourgues and S. Pastor, Adv. High Energy Phys. **2012**, 608515 (2012); H. Ishida, M. Kusakabe, and H. Okada, Phys. Rev. D **90**, 083519 (2014); M. Archidiacono, S. Hannestad, R. S. Hansen, and T. Tram, J. Cosmol. Astropart. Phys. 08 (2016) 067; K. N. Abazajian and M. Kaplinghat, Annu. Rev. Nucl. Part. Sci. **66**, 401 (2016).

[22] Although Majorana neutrino masses have often been regarded in the past as more generic, many ultraviolet extensions of the SM lead to Dirac neutrinos instead. See, e.g., K. Dick, M. Lindner, M. Ratz, and D. Wright, Phys. Rev. Lett. **84**, 4039 (2000); D. A. Demir, L. L. Everett, and P. Langacker, Phys. Rev. Lett. **100**, 091804 (2008); Y. Grossman and D. J. Robinson, J. High Energy Phys. 01 (2011) 132; N. Memenga, W. Rodejohann, and H. Zhang, Phys. Rev. D **87**, 053021 (2013); A. Aranda, C. Bonilla, S. Morisi, E. Peinado, and J. W. F. Valle, Phys. Rev. D **89**, 033001 (2014); E. Ma and O. Popov, Phys. Lett. B **764**, 142 (2017); A. Dasgupta, S. K. Kang, and O. Popov, arXiv:1903.12558 and references therein.

[23] B. W. Lee and R. E. Shrock, Phys. Rev. D **16**, 1444 (1977). The theorem proved in this work is that the necessary and sufficient condition for the weak leptonic neutral current to be diagonal in mass eigenstates is that leptons of the same chirality and charge must have the same weak T and T_3.

[24] J. Schechter and J. W. F. Valle, Phys. Rev. D **22**, 2227 (1980).

[25] R. E. Shrock, Phys. Lett. B **96**, 159 (1980).

[26] R. E. Shrock, Phys. Rev. D **24**, 1232 (1981).

[27] R. E. Shrock, Phys. Rev. D **24**, 1275 (1981).

[28] R. E. Shrock, in *Proceedings of the 1980 VPI Conference on Weak Interactions as Probes of Unification, AIP Conf. Proc. 72*, Particles and Fields Subseries No. 23 (AIP, New York, 1981), p. 368.

[29] J. Deutsch, M. LeBrun, and R. Prieels, Nucl. Phys. **A518**, 149 (1990).

[30] A. Kusenko, Phys. Rep. **481**, 1 (2009); A. Boyarsky, O. Ruchayskiy, and M. Shaposhnikov, Annu. Rev. Nucl. Part. Sci. **59**, 191 (2009).

[31] J. C. Helo, S. Kovalenko, and I. Schmidt, Nucl. Phys. **B853**, 80 (2011); Phys. Rev. D **84**, 053008 (2011).

[32] A. Abada, D. Das, A. M. Teixeira, A. Vicente, and C. Weiland, J. High Energy Phys. 02 (2013) 048; A. Abada, A. M. Teixeira, A. Vicente, and C. Weiland, J. High Energy Phys. 02 (2014) 091.

[33] A. de Gouvêa and A. Kobach, Phys. Rev. D **93**, 033005 (2015).

[34] S. Alekhin *et al.*, Rep. Prog. Phys. **79**, 124201 (2016); M. Drewes *et al.*, J. Cosmol. Astropart. Phys. 01 (2017) 025; K. N. Abazajian, Phys. Rep. **711–712**, 1 (2017); A. Boyarsky, M. Drewes, T. Lasserre, S. Mertens, and O. Ruchayskiy, Prog. Part. Nucl. Phys. **104**, 1 (2019).

[35] B. Batell, T. Han, D. McKeen, and B. Es Haghi, Phys. Rev. D **97**, 075016 (2018).

[36] J. C. Hardy and I. S. Towner, Nucl. Phys. **A254**, 221 (1975).

[37] J. C. Hardy, I. S. Towner, V. T. Kozlowsky, E. Hagberg, and H. Schmeing, Nucl. Phys. **A509**, 429 (1990).

[38] A. Czarnecki, W. J. Marciano, and A. Sirlin, Phys. Rev. D **70**, 093006 (2004).

[39] J. C. Hardy and I. S. Towner, Phys. Rev. C **71**, 055501 (2005); **79**, 055502 (2009).

[40] W. J. Marciano and A. Sirlin, Phys. Rev. Lett. **96**, 032002 (2006).

[41] J. C. Hardy and I. S. Towner, Phys. Rev. C **91**, 025501 (2015).

[42] J. C. Hardy and I. S. Towner, arXiv:1807.01146.

[43] C. Y. Seng, M. Gorchtein, H. H. Patel, and M. Ramsey-Musolf, Phys. Rev. Lett. **121**, 241804 (2018); C. Y. Seng, M. Gorchtein, and M. Ramsey-Musolf, Phys. Rev. D **100**, 013001 (2019).

[44] See talks by J. Hardy, W. Marciano, and M. Ramsey-Musolf, Workshop on First-Row CKM Unitarity (Texas A&M University, 2019), https://cyclotron.tamu.edu/ckmuw2019; See also A. Czarnecki, W. J. Marciano, and A. Sirlin, arXiv:1907.06737, which appeared after our present work had been submitted for publication.

[45] Since in the presence of massive neutrinos, $G_{F,\mathrm{app}}$ would differ from G_F, the value of $|V_{ud}|$, denoted $|V_{ud,\mathrm{app}}|$, derived from these fits to $\mathcal{F}t$ values would also differ from the true value [26,28]. However, all that we use here is the relative agreement between the $\mathcal{F}t$ values, as measured by the quoted precision in the determination of $|V_{ud,\mathrm{app}}|$, so we keep the subscript app implicit.

[46] Using their value of $|V_{ud}|$ in [42], together with $|V_{us}| = 0.2243(5)$ from [11] (the $|V_{ub}|^2$ term makes a negligible contribution), Hardy and Towner obtain $|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 0.99939(64)$, in agreement with CKM unitarity. Reference [43] notes that its value of $|V_{ud}|$ leads to tension with CKM unitarity. See also E. Blucher and W. J. Marciano, in [11].

[47] A. Aguilar-Arevalo *et al.* (PIENU Collaboration), Phys. Rev. D **97**, 072012 (2018).

[48] C. Lazzeroni *et al.* (NA62 Collaboration), Phys. Lett. B **772**, 712 (2017); E. Cortina Gill *et al.* (NA62 Collaboration), Phys. Lett. B **778**, 137 (2018).

[49] D. Bryman, R. Dubois, T. Numao, B. Olaniyi, A. Olin, M. S. Dixit, J.-M. Poutissou, and J. A. Macdonald, Phys. Rev. Lett. **50**, 1546 (1983).

[50] R. Abela, M. Daum, G. H. Eaton, R. Frosch, B. Jost, P.-R. Kettle, and E. Steiner, Phys. Lett. **105B**, 263 (1981); R. C. Minehart, K. O. H. Ziock, R. Marshall, W. A. Stephens, M. Daum, B. Jost, and P.-R. Kettle, Phys. Rev. Lett. **52**, 804 (1984).

[51] Y. Asano *et al.*, Phys. Lett. **104B**, 84 (1981); R. S. Hayano *et al.*, Phys. Rev. Lett. **49**, 1305 (1982).

[52] T. Yamazaki, in *Proceedings of the Neutrino-84 Conference, Nordkirchen, Germany* (World Scientific, Singapore, 1984).

[53] M. Daum, P.-R. Kettle, B. Jost, R. M. Marshall, R. C. Minehart, W. A. Stephens, and K. O. H. Ziock, Phys. Rev. D **36**, 2624 (1987).
[54] D. I. Britton *et al.*, Phys. Rev. D **46**, R885 (1992).
[55] D. Bryman and T. Numao, Phys. Rev. D **53**, 558 (1996).
[56] C. Lazzeroni *et al.* (NA62 Collaboration), Phys. Lett. B **719**, 326 (2013).
[57] A. V. Artamonov *et al.* (BNL E949), Phys. Rev. D **91**, 052001 (2015).
[58] A. Aguilar-Arevalo *et al.* (PIENU Collaboration), Phys. Rev. Lett. **115**, 071801 (2015).
[59] C.-S. Park *et al.* (Belle Collaboration), Phys. Rev. D **94**, 012003 (2016).
[60] A. Aguilar-Arevalo *et al.* arXiv:1904.03269.
[61] S. Berman, Phys. Rev. Lett. **1**, 468 (1958); T. Kinoshita, Phys. Rev. Lett. **2**, 477 (1959); T. Kinoshita and A. Sirlin, Phys. Rev. **113**, 1652 (1959).
[62] W. J. Marciano and A. Sirlin, Phys. Rev. Lett. **71**, 3629 (1993).
[63] V. Cirigliano and I. Rosell, Phys. Rev. Lett. **99**, 231801 (2007); J. High Energy Phys. 10 (2007) 005.
[64] Sufficiently soft photons are understood to be included in the experimentally measured $M^+ \to \ell^+ \nu_\ell$ decays, which are thus often labeled as $M^+ \to \ell^+ \nu_\ell(\gamma)$.
[65] D. Bryman, W. J. Marciano, R. Tschirhart, and T. Yamanaka, Annu. Rev. Nucl. Part. Sci. **61**, 331 (2011).
[66] Hadronic structure-dependent bremsstrahlung (SDB) contributions to the radiative corrections are estimated to be negligibly small for $R_{e/\mu}^{(\pi)}$ and hence are not included in the calculations of [62,63]. Bremsstrahlung contributions to $R_{e/\mu}^{(K)}$ are discussed in [48,63].
[67] D. A. Bryman, R. Dubois, J. A. Macdonald, T. Numao, B. Olaniyi, A. Olin, J.-M. Poutissou, and M. S. Dixit, Phys. Rev. D **33**, 1211 (1986).
[68] D. I. Britton *et al.*, Phys. Rev. Lett. **68**, 3000 (1992).
[69] G. Czapek *et al.*, Phys. Rev. Lett. **70**, 17 (1993).
[70] G. Feldman and R. Cousins, Phys. Rev. D **57**, 3873 (1998).
[71] Charge-conjugate decays are understood to be included here and below.
[72] J. P. Alexander *et al.* (CLEO Collaboration), Phys. Rev. D **79**, 052001 (2009).
[73] P. del Amo Sanchez *et al.* (*BABAR* Collaboration), Phys. Rev. D **82**, 091103(R) (2010).
[74] A. Zupanc *et al.* (Belle Collaboration), J. High Energy Phys. 09 (2013) 139.
[75] M. Ablikim *et al.* (BES III Collaboration), Phys. Rev. D **94**, 072004 (2016).
[76] M. Ablikim *et al.* (BES III Collaboration), Phys. Rev. Lett. **122**, 071802 (2019).
[77] C. Yuan and G. Zhao (BES III Collaboration) (private communication).
[78] N. Satoyama *et al.* (Belle Collaboration), Phys. Lett. B **647**, 67 (2007).
[79] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. D **79**, 091101 (2009).
[80] A. Sibidanov *et al.* (Belle Collaboration), Phys. Rev. Lett. **121**, 031801 (2018).
[81] M. Prim (Belle Collaboration), proceedings at Moriond-2019 [82].
[82] Moriond-2019 Workshop, http://moriond.in2p3.fr/2019/EW.
[83] J. P. Lees *et al.* (*BABAR* Collaboration), Phys. Rev. D **88**, 031102 (2013).
[84] K. Hara *et al.* (Belle Collaboration), Phys. Rev. Lett. **110**, 131801 (2013).
[85] B. Kronenbitter *et al.* (Belle Collaboration), Phys. Rev. D **92**, 051102 (2015).
[86] See CKMfitter website, http://ckmfitter.in2p3,fr.
[87] Here, we focus on the effect of ν_4 in these B decays. For recent discussions of possible indications of other new physics in the B system, see, e.g., [82].
[88] W. J. Marciano and A. I. Sanda, Phys. Rev. Lett. **67**, 303 (1977); S. M. Bilenky, S. T. Petcov, and B. Pontecorvo, Phys. Rev. Lett. **67**, 309 (1977); B. W. Lee, S. Pakvasa, R. E. Shrock, and H. Sugawara, Phys. Rev. Lett. **38**, 1230 (1977).
[89] K. Fujikawa and R. E. Shrock, Phys. Rev. Lett. **45**, 963 (1980).
[90] R. E. Shrock and M. Suzuki, Phys. Lett. **110B**, 250 (1982).
[91] S. Cooperstein (CMS Collaboration), in [82]; see also V. Khachatryan *et al.* (CMS Collaboration), J. High Energy Phys. 02 (2017) 135; Phys. Lett. B **793**, 520 (2019); G. Aad (ATLAS Collaboration), Phys. Lett. B **793**, 499 (2019).
[92] B. Bernardi *et al.*, Phys. Lett. **166B**, 479 (1986).