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Decaying Sterile Neutrinos and the Short Baseline Oscillation Anomalies

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The MiniBooNE experiment has observed a significant excess of electron neutrinos in a muon neutrino beam at source–detector distances too short to be compatible with standard neutrino oscillations. The most straightforward explanation for this signal in terms of oscillations between Standard Model neutrinos and a new, sterile, neutrino, is disfavored by null results from experiments looking for muon neutrino disappearance. Here, we discuss the possibility that MiniBooNE data are instead explained by a sterile neutrino that decays quickly back into active neutrinos plus a light boson. The flavor composition of the secondary neutrinos is determined by the sterile neutrino mixing angles, and we show that the data is best explained if the sterile neutrino mixes mostly with electron neutrinos. The preferred range for the mass of the sterile neutrino is between 100 eV and 1 keV. We argue that the model can easily satisfy cosmological constraints because it has the "secret interactions" mechanism built-in. Accommodating in addition to the MiniBooNE anomaly also the LSND, reactor, and gallium anomalies is possible, but in this case the model needs to be extended to avoid cosmological limits.

Many major discoveries in neutrino physics have started out as oddball anomalies that gradually evolved into incontrovertible evidence. In this letter, we entertain the possibility that history is repeating itself in the context of the MiniBooNE anomaly. From 2002 to 2019, the MiniBooNE experiment has been searching for electron neutrinos (ν_e) appearing in a muon neutrino (ν_μ) beam [1–3],[1] and has found a corresponding signal at 4.8σ statistical significance. For some time, the simplest explanation for this signal appeared to be the existence of a fourth neutrino species ν_s, called "sterile neutrino" because it would not couple to any of the Standard Model interactions, but would communicate with the Standard Model only via neutrino mixing. If ν_s has small but non-zero mixing with both ν_e and ν_μ and if the corresponding mostly sterile neutrino mass eigenstate ν_4 is somewhat heavier ($\sim$ 1 eV) than the Standard Model neutrinos, the MiniBooNE signal could be explained. This explanation would also be consistent with a similar 3.8σ anomaly from the earlier LSND experiment [4], and with several reported hints for anomalous disappearance of electron neutrinos in reactor experiments [5, 6] and in experiments using intense radioactive sources [7, 8].[2] However, the sterile neutrino parameter space consistent with MiniBooNE and these other anomalies is in severe tension with the non-observation of anomalous ν_μ disappearance [9–19], unless several additional new physics effects are invoked concomitantly [20].

In this letter, we propose a different explanation for the MiniBooNE anomaly, and possibly also for the LSND, reactor, and gallium anomalies. In particular, we consider a sterile neutrino that rapidly decays back into Standard Model ("active") neutrinos [21, 22]. The MiniBooNE excess is then interpreted as coming from these decay products. We will see that this scenario requires only very small mixing between ν_s and ν_μ, thus avoiding the strong ν_μ disappearance constraints. It also requires somewhat larger mixing between ν_s and ν_e, in line with the hints from reactor and radioactive source experiments. Finally, we will argue that decaying sterile neutrinos may avoid cosmological constraints because the model automatically endows sterile neutrinos with self-interactions ("secret interactions" [23, 24]).

Decaying Sterile Neutrino Formalism. We extend the standard model by a sterile neutrino ν_s (a Dirac fermion) and a singlet scalar ϕ. The relevant interaction and mass terms in the Lagrangian of the model are

$$\mathcal{L} \supset -g\,\bar\nu_s \nu_s \phi - \sum_{a=e,\mu,\tau,s} m_{\alpha\beta}\,\bar\nu_\alpha \nu_\beta\,. \qquad (1)$$

The neutrino flavor eigenstates ν_α are linear combinations of the mass eigenstates ν_j ($j = 1, 2, 3, 4$) according to the relation $\nu_\alpha = U_{\alpha j}\nu_j$, where U is the unitary 4×4 leptonic mixing matrix. The first term in eq. (1) can thus be rewritten as

$$-g\,\bar\nu_F \nu_F \phi - g\,|U_{s4}|^2 \bar\nu_4 \nu_4 \phi - (g\,U^*_{s4} \bar\nu_4 \nu_F \phi + h.c.)\,, \qquad (2)$$

with $\nu_F \equiv \sum_{i=1}^{3} U_{si}\nu_i$. We assume initially that the fourth, mostly sterile, mass eigenstate $\nu_4 \simeq \nu_s$ has a mass

[1] Here and in the following, when we say neutrino we mean also the corresponding anti-neutrinos.

[2] The latter class of experiments is usually referred to as "gallium experiments", based on the active component of their target material.

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m_4 between $\mathcal{O}$(eV) and $\mathcal{O}$(100 keV), and that the mass of ϕ is of the same order, but smaller. The last term in eq. (2) will then induce $\nu_4 \to \nu_F + \phi$ decays, while the first term is responsible for $\phi \to \nu_F + \bar{\nu}_F$ decays. When these decays occur in a neutrino beam, they will produce lower-energy neutrinos at the expense of higher-energy ones, and they may also alter the flavor structure of the beam. In particular, they can produce excess low-energy ν_e in a ν_μ beam, as suggested by the MiniBooNE anomaly.

The phenomenology of the model depends mainly on five new parameters. Besides m_4 and m_ϕ, these are the coupling g and the mixings $|U_{e4}|^2$, $|U_{\mu 4}|^2$ between ν_4 and ν_e, ν_μ. We will assume the mixing with ν_τ to be zero and neglect the complex phases, as these parameters do not play an important role in explaining the MiniBooNE excess. For practical purposes, it is convenient to quote $m_4\Gamma_4$ instead of g, as $m_4\Gamma_4$ appears directly in the laboratory frame decay length $E/(m_4\Gamma_4)$. Also, it is convenient to use the ratio m_ϕ/m_4 instead of just m_ϕ because the ratio measures more directly the kinematic suppression in ν_4 decays.

The evolution in energy E and time t of a neutrino beam in our model can be described by a neutrino density matrix $\hat{\rho}_\nu(E, x)$ (a 4×4 matrix in flavor space), the corresponding antineutrino density matrix $\hat{\bar{\rho}}_\nu(E, x)$, and the scalar density function $\rho_\phi(E, t)$. The evolution equations are [25, 26],

$$\frac{d\hat{\rho}_\nu(E,t)}{dt} = -i[\hat{H}, \hat{\rho}_\nu] - \frac{1}{2}\big\{ \tfrac{m_4}{E}\hat{\Gamma}, \rho \big\} + \mathcal{R}_\nu[\hat{\rho}_\nu, \rho_\phi, E, t] \tag{3}$$

$$\frac{d\rho_\phi(E,t)}{dt} = -\tfrac{m_\phi}{E}\Gamma_\phi \rho_\phi + \mathcal{R}_\phi[\hat{\rho}_\nu, E, t] \tag{4}$$

where $\hat{H} = \frac{1}{2E}\,\mathrm{diag}(0, \Delta m_{21}^2, \Delta m_{31}^2, \Delta m_{41}^2)$ is the standard neutrino oscillation Hamiltonian, written here in the mass basis, and $\hat{\Gamma} = \Gamma_4 \hat{\Pi}_4$ is the decay term, which contains the projection operator $\hat{\Pi}_4 = |\nu_4\rangle\langle\nu_4|$ onto the fourth, mostly sterile, mass eigenstate as well as the decay width Γ_4 of ν_4 in its rest frame. Similarly, Γ_ϕ is the rest frame decay width of ϕ. The functional $\mathcal{R}_\nu[\hat{\rho}_\nu, \rho_\phi, E, t]$ describes the appearance of the daughter neutrinos from ν_4 and ϕ decay. Neglecting the masses of ν_1, ν_2, and ν_3, it is given by

$$\mathcal{R}_\nu[\hat{\rho}_\nu, \rho_\phi, E, t] = \hat{\Pi}_F \int_{\frac{E}{1-x_{\phi 4}^2}}^{\infty} dE_4 \sum_k \hat{\rho}_{\nu,44}(E_4, t)\, \frac{d\Gamma^{\mathrm{lab}}(\nu_4 \to \nu_k \phi)}{dE_k} \;+\; \hat{\Pi}_F \sum_{k,j} \int_E^\infty dE_\phi\, \rho_\phi(E_\phi, t) \frac{d\Gamma^{\mathrm{lab}}(\phi \to \nu_k \bar{\nu}_j)}{dE_k}\,, \tag{5}$$

where $d\Gamma^{\mathrm{lab}}(X \to Y)/dE_k$ are the differential decay widths for the various decays $X \to Y$ in the lab frame, and $x_{\phi 4} \equiv m_\phi/m_4$. The projection operator

$$\hat{\Pi}_F = \frac{|\nu_F\rangle\langle\nu_F|}{|\langle\nu_F|\nu_F\rangle|^2} = \sum_{i,j=1}^{3} \frac{U_{si}^* U_{sj}}{\sum_k |U_{sk}|^2} |\nu_i\rangle\langle\nu_j| \tag{6}$$

isolates the specific combination of mass eigenstates that appears in ν_4 and ϕ decays, and the integrals run over all parent energies E_4, E_ϕ that lead to daughter neutrinos of energy E. Analogously, $\mathcal{R}_\phi[\hat{\rho}_\nu, E, t]$ describes the appearance of scalars from ν_4 decay:

$$\mathcal{R}_\phi[\hat{\rho}_\nu, E, t] = \int_E^{E/x_{\phi 4}^2} dE_4 \sum_k \left[\hat{\rho}_{\nu,44}(E_4, t)\, \frac{d\Gamma^{\mathrm{lab}}(\nu_4 \to \nu_k \phi)}{dE_\phi} + \hat{\bar{\rho}}_{\nu,44}(E_4, t) \frac{d\Gamma^{\mathrm{lab}}(\bar{\nu}_4 \to \bar{\nu}_k \phi)}{dE_\phi} \right]. \tag{7}$$

Analytic expressions for the decay widths and the $\nu_\mu \to \nu_e$ transition probability are given in the Appendix.

Fit to MiniBooNE data. To compare the predictions of the decaying sterile neutrino scenario to MiniBooNE data, we evolve the unoscillated beam following the formulas given above. We then follow the fitting procedure recommended by the MiniBooNE collaboration (see the data releases accompanying refs. [1, 3]), but go beyond it by accounting for the impact of ν_μ and ν_e disappearance on the signal and background normalization (see Appendix for details).

Illustrative results are shown in fig. 1, where we have chosen parameter values that give an optimal fit to MiniBooNE data while being consistent with null results from other oscillation experiments, as well as non-oscillation constraints. At $m_4\Gamma_4 = 2.1\,\mathrm{eV}^2$, most ν_4 will have decayed before reaching the detector. The value $m_\phi/m_4 = 0.82$ implies mild phase space suppression in ν_4 decays, which tends to shift the ν_e spectrum to lower energies, in excellent agreement with the data. Compared to models

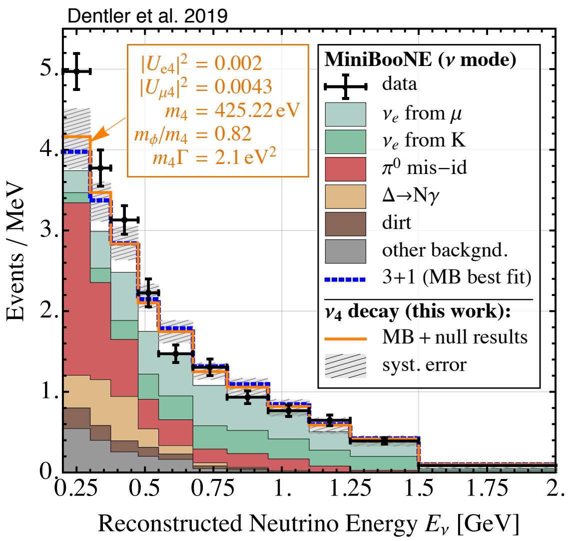

FIG. 1. Comparison of MiniBooNE data [3] to the predictions of the neutrino oscillation + decay scenario discussed in this letter. We show the expected spectrum at the point which optimally fits MiniBooNE data, while being consistent with all null results (orange histogram with systematic error band; parameters given in the plot). We also show the MiniBooNE-only best point for $3 + 1$ oscillations without decay (blue dotted histogram, parameter values $\Delta m^2_{41} = 0.13\,\text{eV}^2$, $|U_{e4}|^2 = 0.024$, $|U_{\mu4}|^2 = 0.63$).

with massless ϕ [21, 22], our scenario also has the advantage that it allows $\phi \to \nu_F \bar{\nu}_F$ decays, further boosting the ν_e flux at low energies. It is therefore favored compared to the $m_\phi = 0$ case at more than 99% confidence level. The fit in our model is better than in oscillation-only scenarios (blue dotted histogram in fig. 1) [19], which by themselves already offer an excellent fit as long as only MiniBooNE data are considered (MiniBooNE quotes a χ^2 per degree of freedom of $9.9/6.7$ [3]). Our model, however, is also consistent with all constraints. Notably, it reproduces the angular distribution of the neutrino interaction products in MiniBooNE because it predicts an actual flux of electron neutrinos instead of attempting to mimic the signal with other particles [27–32].

Constraints. We now discuss the various constraints that an explanation of the MiniBooNE anomaly in terms of decaying sterile neutrinos has to respect. The most relevant constraints are also summarized in figs. 2 and 3.

(1) Oscillation null results. Putting MiniBooNE into context with other ν_e appearance searches, we show in fig. 2 two slices through the 5-dimensional parameter space of the decaying sterile neutrino model along the plane spanned by $|U_{e4}|^2$ and $|U_{\mu4}|^2$. To produce this figure, we have used fitting codes from refs. [9, 12, 19] (based partly on refs. [33–35]). We see that most of the parameter region preferred by MiniBooNE is well compatible with the KARMEN short-baseline oscillation search [36] and with the OPERA long-baseline experiment [37]. We have checked that the limits from ICARUS [38–40] and E776 [41] are significantly weaker.

Potentially relevant constraints on $|U_{e4}|^2$ could arise from searches for ν_e disappearance using reactor neutrinos, neutrinos from intense radioactive source, solar neutrinos, atmospheric neutrinos, and neutrinos from pion decay at rest [19]. However, as large $|U_{e4}|^2$ implies that the invisible ν_4 quickly decay back to ν_e, we expect no net reduction of the ν_e flux in these searches. In fact, one may even expect an increase due to the ν_e from ϕ decay. Since ν_e produced in decays will have lower energies than their parent particles, some spectral distortions are expected. As many ν_e disappearance limits are based on total rate measurements (atmospheric, π decay at rest, much of the solar neutrino data) these limits will be significantly weakened compared to the oscillation-only scenario.

Also constraints from ν_μ disappearance, which are prohibitive in non-interacting sterile neutrino models [17, 19, 42], are irrelevant here. First, at $m_4 \gg$ eV, these constraints are much weaker than at $m_4 \sim 1\,\text{eV}$, the favored mass range in oscillation-only models. This is because analyses at large m_4 will only see an overall ν_μ flux deficit, but no spectral features. The best available constraint at $m_4 \gg$ eV comes from MINOS+ [43] and is at the level of $|U_{\mu4}|^2 \lesssim 0.02$ (see also ref. [44]). Second, in pure oscillation scenarios the number of excess events in MiniBooNE and LSND is proportional to $|U_{e4}|^2|U_{\mu4}|^2$, while in our scenario it is proportional only to $|U_{\mu4}|^2$ as long as $|U_{e4}|^2 \gg |U_{\mu4}|^2$. Therefore, significantly smaller values of $|U_{\mu4}|^2$ are viable.

We can already see from fig. 2 that MiniBooNE is also compatible with LSND and with the $|U_{e4}|^2$ range preferred by the reactor anomaly, but only in a parameter region that would unacceptably reduce free-streaming of active neutrinos in the early Universe. We will see below that this tension can be avoided in extensions of the model.

(2) Beta decay spectra (purple regions in fig. 3 and black dashed lines in fig. 2). Direct searches for sterile neutrinos looking for anomalous features in beta decay spectra [45–48] suggest that $\mathcal{O}(0.001 - 0.01)$ mixings between active and sterile neutrinos – as required by MiniBooNE – are allowed for $m_4 \lesssim$ few keV.

(3) Neutrinoless double beta decay. If neutrinos are Majorana particles, the non-observation so far of neutrinoless double beta decay requires $m_4|U_{e4}|^2 \lesssim 0.2\,\text{eV}$ [49]. This is the reason we always focus on Dirac neutrinos in this work.

(4) N_{eff}, a measure for the *energy density of relativistic particles in the early Universe* (green region in fig. 3). The measured value of N_{eff} is very close to the SM value of ~ 3 both at the BBN and recombination epochs [50, 51]. Naively, one might expect that this observation precludes the existence of a fourth neutrino species with $m_4 \lesssim$ MeV. In our model, however, the N_{eff} constraint is

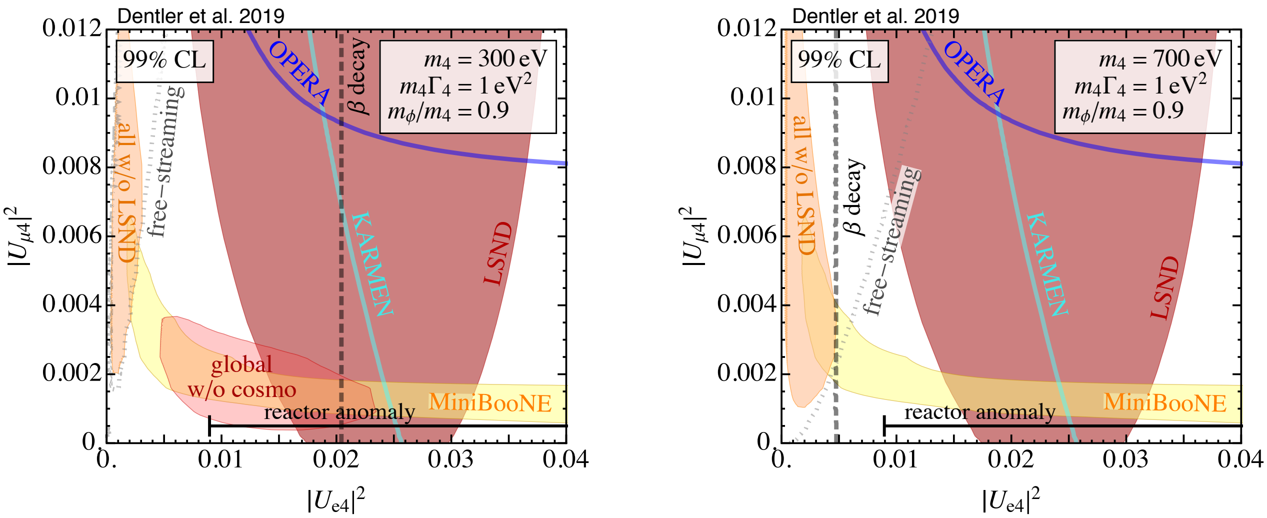

FIG. 2. Allowed values of the squared mixing matrix elements $|U_{e4}|^2$ and $|U_{\mu4}|^2$ (measuring the mixing of ν_s with ν_e and ν_μ, respectively) in the decaying sterile neutrino scenario. We show two representative slices through the 5-dimensional 99% confidence regions. Our fits include MiniBooNE, OPERA, ICARUS, E776, and KARMEN data, as well as constraints from nuclear beta decay spectra and from the requirement of neutrino free-streaming in the early Universe. For the null results, the region to the right of the curves is excluded. We also show, as a black rule at the bottom of the plot, the $|U_{e4}|^2$ range preferred by the reactor neutrino anomaly. Constraints on ν_μ disappearance are significantly weaker here than in the $3+1$ scenario without decay, and are hence not shown. We also do not show a fit including both LSND and cosmology as the goodness of fit would be very poor. Note that the global combinations are sensitive to five degrees of freedom, namely m_4, $|U_{e4}|^2$, $U_{\mu4}^2$, $m_4\Gamma_4$, and m_ϕ/m_4; oscillation experiments are sensitive only to the last four of these; beta decay spectra and free-streaming depend on two degrees of freedom (m_4 and $|U_{e4}|^2$); reactor experiments depend only on $|U_{e4}|^2$.

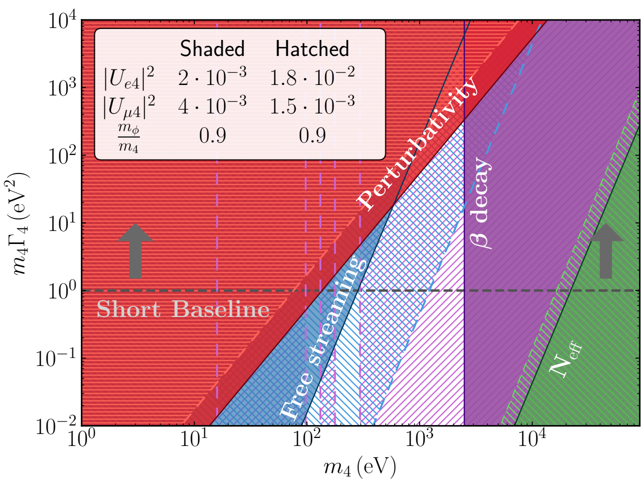

FIG. 3. Non-oscillation constraints on decaying sterile neutrinos for parameters favored by the global fit without LSND (shaded), and by the global fit without the free-streaming constraint (hatched).

avoided by the "secret interactions" mechanism [23, 24]: a small abundance of ν_s produced either via oscillations or at the end of inflation, generates a large, temperature-dependent potential $V_{\rm eff} \propto g^2 T$ for ν_s. This potential suppresses the ν_s–ν_a mixing angle in matter, θ_m, by a factor $\sqrt{\Delta m^2/(E V_{\rm eff})}$ until the temperature drops low enough for this factor to become $\gtrsim 1$. For the parameter range that the short-baseline anomalies are pointing to, this can easily be postponed to late times ($T \ll$ MeV), when the neutrino sector has decoupled from the photons. Consequently, when ν_s are eventually produced, they are produced at the expense of active neutrinos, so $N_{\rm eff}$ does not change any more and constraints are automatically satisfied. More quantitatively, $N_{\rm eff}$ constraints are avoided when

$$(m_4\Gamma_4)^{\rm eff} \gtrsim 2\times10^{-14}\,\mathrm{eV}^2 \left(\frac{m_4}{\mathrm{eV}}\right)^4, \qquad (8)$$

where we have defined

$$(m_4\Gamma_4)^{\rm eff} \equiv \frac{m_4\Gamma_4}{|U_{s4}|^2(|U_{e4}|^2+|U_{\mu4}|^2)\left(1-\frac{m_\phi^2}{m_4^2}\right)^2}. \qquad (9)$$

This constraint can be easily satisfied in the mass range allowed by beta decay limits.

(5) $\sum m_\nu$*, the sum of neutrino masses.* Massive neutrinos affect the CMB as well as structure formation, and this has for instance allowed the Planck collaboration to set a limit $\sum m_\nu \lesssim 0.12\,\mathrm{eV}$ [51]. In our model, this constraint is easily satisfied because in the interesting parameter range with $m_4 \gg 1\,\mathrm{eV}$ and $m_4\Gamma_4 \gtrsim 1\,\mathrm{eV}^2$, any ν_4 that are produced in the early Universe will have decayed via $\nu_4 \to \nu_{1,2,3} + (\phi \to \nu_{1,2,3}\bar{\nu}_{1,2,3})$ long before recombination and the onset of structure formation.

(6) Neutrino Free-Streaming (blue region in fig. 3 and gray dotted lines in fig. 2). Via the mixing with ν_s, also the light neutrino mass eigenstates $\nu_{1,2,3}$ feel ϕ-mediated interactions and are therefore not fully free-streaming. This may put the model in tension with CMB observations, which require that neutrinos should free-stream from about redshift 10^5 onwards [52, 53].[3] This sets the limit

$$(m_4\Gamma_4)^{\text{eff}} \lesssim 4\times 10^{-10}\,\text{eV}^2 \left(\frac{m_4}{\text{eV}}\right)^4 \left(\frac{0.1}{|U_{s1}|}\right)^4 x_{\phi 4}^2 \,. \quad (10)$$

Note that in fig. 2, this constraint is present even for very small mixings. This is because, at fixed $m_4\Gamma_4$, small mixings need to be compensated by a large coupling g, strengthening the free streaming constraint.

Depending on the mixing angles, the free-streaming limit may be a problem in the m_4 range of interest to us. It could be substantially weakened, however, in extensions of the model, see for instance refs. [59–62]. A particularly simple possibility is to add extra species of light free-streaming particles (for instance extra sterile neutrinos or dark photons) that are produced at the expense of the neutrino sector after neutrino decoupling and compensate for the lack of free-streaming in active neutrinos.

(6) SN 1987A. The fact that neutrinos from supernova 1987A could be observed at Earth without being absorbed through scattering on the cosmic neutrino background constrains neutrino self-interactions [63]. We have checked that, due to mixing suppression, these constraints are avoided in our scenario. Note that supernova cooling, which is sensitive to non-interacting sterile neutrinos, does not constrain our model as ν_4 and ϕ quickly decay to lighter neutrinos that remain trapped in the supernova core.

(7) Decays of SM neutrinos. We have checked that decays of the form $\nu_{2,3} \to \bar{\nu}_1 + 2\nu_1$, mediated by an off-shell ϕ, are always sufficiently rare to be consistent with solar neutrino constraints [64, 65]. Note, however, that we predict the cosmic neutrino background today to consist exclusively of ν_1 or ν_3, for normal and inverted neutrino mass ordering, respectively.

(8) Perturbativity (red region in fig. 3). Requiring that the ν_s–ϕ coupling constant g in eqs. (1) and (2) is $< \sqrt{4\pi}$ imposes the bound

$$(m_4\Gamma_4)^{\text{eff}} \lesssim 0.25\,\text{eV}^2 \left(\frac{m_4}{\text{eV}}\right)^2 . \quad (11)$$

Similarly to the free-streaming bound, this constraint applies even for very small mixing when $m_4\Gamma_4$ is fixed. This bound restricts m_4 in our model to be $\gtrsim 100\,\text{eV}$ for $m_4\Gamma_4$ values large enough to explain the MiniBooNE anomaly.

In summary, the sterile neutrino mass range to explain the MiniBooNE anomaly is between 100 eV and 2.5 keV.

The LSND and reactor anomalies. As shown in fig. 2, decaying sterile neutrinos can simultaneously fit the MiniBooNE and LSND anomalies, but only if cosmological neutrino free-streaming constraints can be avoided (see discussion under point *(6)* above for possible scenarios). Quantitatively, a parameter goodness-of-fit test [66] reveals that LSND is incompatible with the rest of the data at the 4.7σ level if free-streaming constraints hold. If the free-streaming problem is solved by other means, this reduces to 2.1σ, implying consistency. The best fit to all data including LSND, but excluding free-streaming is found at $m_4 = 97\,\text{eV}$, $|U_{e4}|^2 = 0.018$, $|U_{\mu 4}|^2 = 0.0015$, $m_4\Gamma_4 = 0.87\,\text{eV}^2$, $m_\phi/m_4 = 0.89$.

Interestingly, at this value of $|U_{e4}|^2$, the model can also explain the flux deficit observed in reactor and gallium experiments [5–8, 19, 67]. We test our model against reactor data by comparing to Daya Bay's generic flux-weighted cross section [68]. To estimate the viable parameter space we perform a chi-square-test using the covariance matrix given in the same reference. In addition we introduce a 2.4% systematic flux normalization error corresponding to the theoretical uncertainty, in accordance with fig. 28 of ref. [68]. The $|U_{e4}|^2$ region preferred by reactor experiments is included in fig. 2, and a comparison of the reactor neutrino spectrum to our model prediction is shown in fig. 4.

Conclusions. In summary, we have shown that scenarios in which the SM is extended by a sterile neutrino that has a decay mode to active neutrinos can well explain the MiniBooNE anomaly without violating any constraints. An explanation of the LSND and reactor/gallium anomalies is possible if the model is extended to avoid constraints on neutrino free-streaming in the early Universe. The preferred mass of the sterile neutrino is of order few hundred eV.

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[3] It is noteworthy, though, that some cosmological fits have actually found a *preference* for neutrino self-interactions [52, 54–58] that could be accommodated in our model.

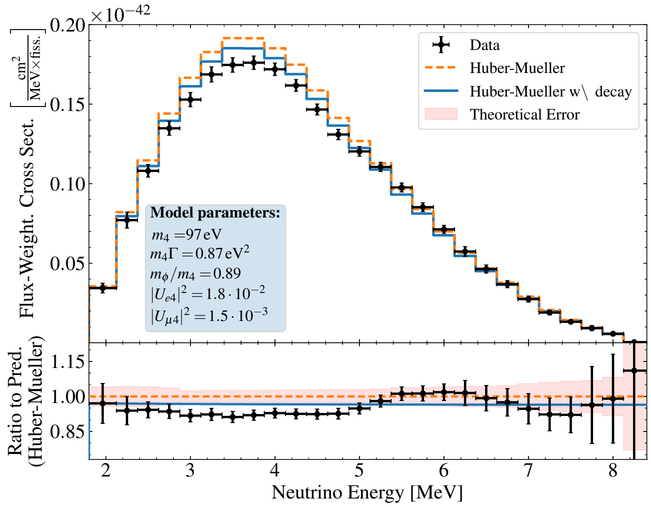

FIG. 4. Comparison of the reactor antineutrino spectrum predicted in the decaying sterile neutrino scenario discussed in this letter (blue) to the standard Huber–Mueller prediction (orange-dashed) [69, 70] and to Daya Bay data (black data points with error bars) [68]. For model parameters motivated by the MiniBooNE and LSND anomalies, a flux deficit consistent with the reactor anomaly can be accommodated. (See text for details, and for a discussion of how possible cosmological constraints can be avoided.)

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Appendix A: Impact of Oscillations on the Background Prediction in MiniBooNE

In this appendix, we briefly discuss our fit to MiniBooNE data, and in what ways it differs from the collaborations' fit as described in the supplemental material to ref. [1], and using the data released with ref. [3]. In particular, we consider the following three effects, which are relevant in a fit to a 3+1 scenario, but are not encountered in a 2-flavor fit.

1. **Normalization of the $\nu_\mu \to \nu_e$ oscillation signal.** To predict the number of expected ν_e events from $\nu_\mu \to \nu_e$ oscillations for a given set of oscillation parameters, the initial ν_μ flux must be known. It is obtained in situ using MiniBooNE's own sample of ν_μ events. Note, however, that in a 3+1 model, the measured ν_μ flux will be reduced by an amount $\sim |U_{\mu 4}|^2$ due to $\nu_\mu \to \nu_s$ oscillations. (This effect is unimportant in a 2-flavor model, where the deficit is only of order $\sin^2 2\theta_{\mu e}$, where $\theta_{\mu e}$ is the effective 2-flavor mixing angle.) We account for this effect by first computing the expected ν_e signal based on the unoscillated MiniBooNE flux, and then diving it by the ν_μ survival probability in each bin.

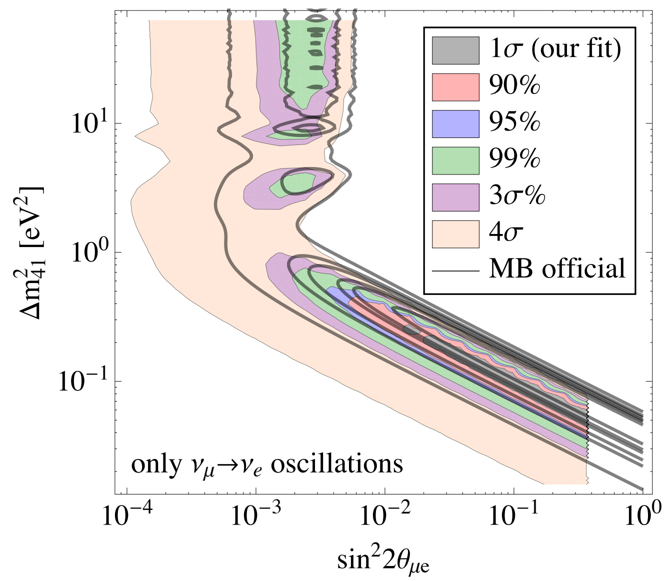
   The impact of this change in normalization is illustrated in the top panels of fig. 5. The colored region in panel (a) of this figure shows our reproduction of the official MiniBooNE fit, which is shown as black contours. In panel (b), we have included the change in normalization for the signal.

2. **Oscillations of the ν_e backgrounds.** Part of the MiniBooNE background is constituted by the intrinsic ν_e contamination in the beam. In a 2-flavor fit, this contribution to the total event rate is only modified by a factor of order $\sin^2 2\theta_{\mu e}$, but in the full 3+1 framework, it is reduced by a factor of order $|U_{e4}|^2$ instead. The impact of this modification to the background sample is shown in fig. 5 (c).

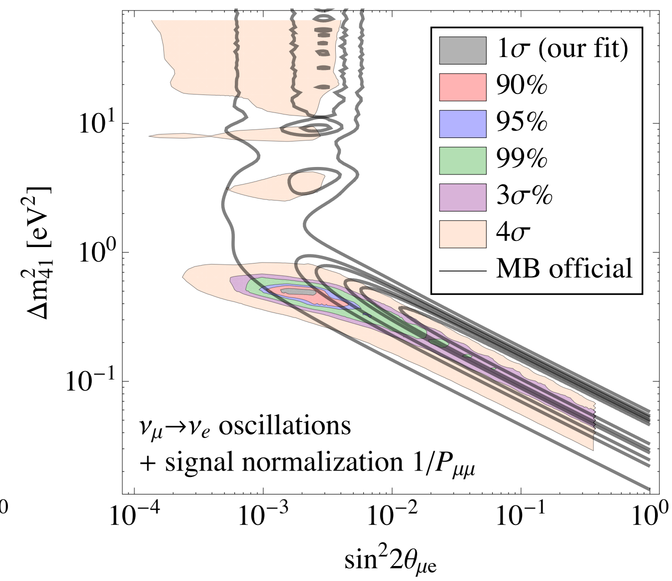
3. **Oscillations of the ν_μ sample.** The fit described in the supplemental material to ref. [1] which we are following includes also MiniBooNE's sample of ν_μ events. This is necessary to properly account for systematic uncertainties which are correlated between the two samples. But of course, in a 3+1 scenario, the ν_μ sample suffers from ν_μ disappearance into ν_s, proportional to $|U_{\mu 4}|^2$. (Once again, in a 2-flavor model, only a much smaller fraction $\propto \sin^2 2\theta_{\mu e}$ will disappear, which is usually negligible.) The impact of including ν_μ disappearance is shown in panel (d) of fig. 5.

We see that including the effect of 3+1 oscillations on the normalization in the control regions and on the background prediction reduces the significance of the MiniBooNE anomaly, though it remains above 3σ. These effects are thus unable to fully explain the MiniBooNE anomaly, but they could well be part of an "Altarelli cocktail" of several effects conspiring to lead to the large observed excess [71].

Let us finally mention one caveat with the above corrections to the MiniBooNE fit. Namely, we can only apply the corrections at the level of reconstructed events as the mapping between true and reconstructed neutrino energies is not publicly available for muon neutrinos. This means we have to assume that the reconstructed neutrino energy is a faithful representation of the true neutrino energy. While this is true for quasi-elastic scattering events which constitute the majority of events, it is not the case for other event categories. For instance, a neutrino–nucleon interaction may create an extra pion, and if this pion is reabsorbed as it propagates out of the nucleus, the event will be misinterpreted as a quasi-elastic interaction, and the kinematic reconstruction of the neutrino energy based on the observed charged lepton energy and direction will fail.



(a)



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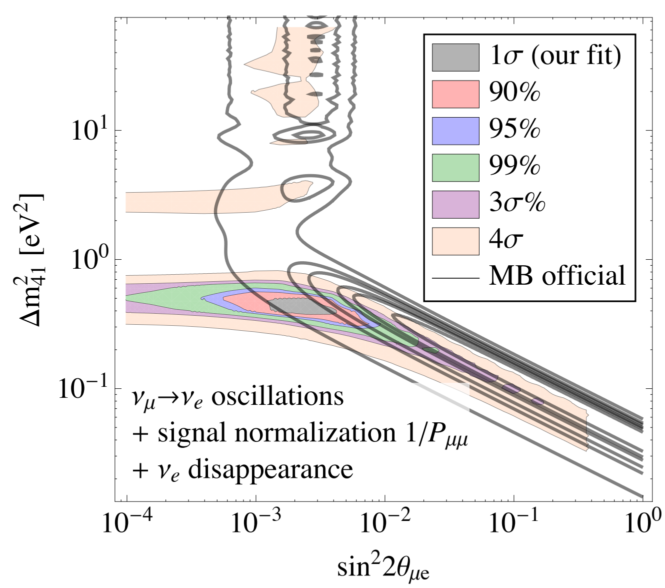

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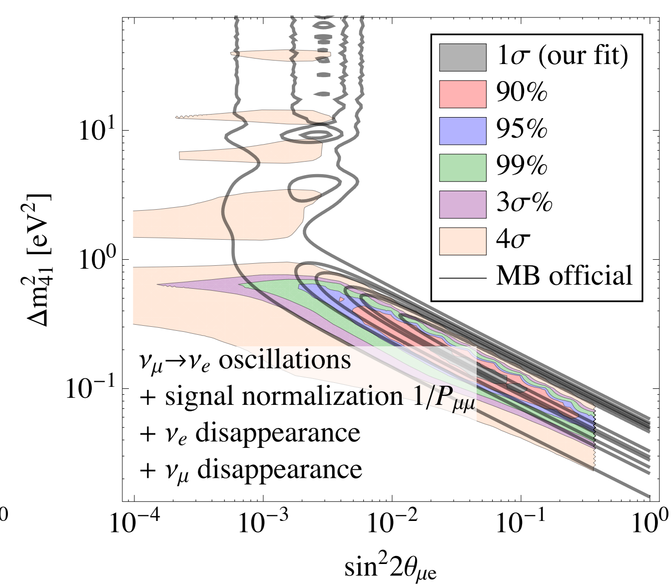

(d)

FIG. 5. Impact of oscillations in the background and control regions on the MiniBooNE fit in a simple 3+1 model (oscillations only, no decay). All panels show Δm^2_{41} vs. the effective 2-flavor mixing angle $\sin^2 2\theta_{\mu e}$, which in a 3+1 scenario is given by $4|U_{e4}|^2|U_{\mu 4}|^2$. Panel (a) shows our reproduction (colored regions) of the official MiniBooNE fit (black contours), based on the instructions given in the supplemental material to ref. [1] and using the data released with ref. [3]. In panel (b), we include in addition the impact of $\nu_\mu \to \nu_s$ disappearance on the normalization of the signal in each bin. The colored contours in panel (c) include on top of this the effect of ν_s disappearance on the intrinsic ν_e contamination in the beam. Panel (d) finally shows the additional impact of ν_μ disappearance on the sample of ν_μ events that is included in the fit along with the ν_e sample. In all panels, we show projections of the three-dimensional parameter space spanned by Δm^2_{41}, $|U_{e4}|^2$, and $U_{\mu 4}|^2$ onto the Δm^2_{41}–$\sin^2 2\theta_{\mu e}$ plane, imposing the constraint $|U_{e4}|^2 < 0.2$ due to bounds from reactor neutrino experiments.

Appendix B: Decay widths and transition probability

Decay rates

Based on the interaction terms from eq. (2), we can compute the differential decay rates of the heavy neutrino ν_4 and of the scalar ϕ. In the massless light neutrino limit, we obtain for the ν_4 decay width in the lab

$$\frac{1}{\frac{m_4}{E_4}\Gamma_4}\frac{d\Gamma^{\rm lab}(\nu_4\to\nu_j\phi)}{dE_j}=\frac{|U_{sj}|^2}{\sum_{k=1}^3|U_{sk}|^2}\frac{E_j}{(1-x_{\phi4}^2)^2E_4^2}\,, \tag{12}$$

$$\sum_j\frac{1}{\frac{m_4}{E_4}\Gamma_4}\frac{d\Gamma^{\rm lab}(\nu_4\to\nu_j\phi)}{dE_\phi}=\frac{1}{1-x_{\phi4}^2}\frac{1}{E_4}\,. \tag{13}$$

In these expressions,

$$\Gamma_4=\frac{g^2}{16\pi}m_4(1-x_{\phi4}^2)^2\sum_{j=1}^3|U_{s4}^*U_{sj}|^2 \tag{14}$$

is the total rest frame decay width of ν_4, $x_{\phi4}\equiv m_\phi/m_4$ is the ratio of scalar and neutrino masses, and E_j, E_ϕ are the daughter neutrino and scalar energies, respectively. In the ν_4 rest frame, E_j is restricted to the interval $[0, m_4(1-x_{\phi4}^2)]$.

The lab frame decay rate of the scalar ϕ is

$$\sum_{i,j}\frac{1}{\frac{m_\phi}{E_\phi}\Gamma_\phi}\frac{d\Gamma^{\rm lab}(\phi\to\nu_i\bar\nu_j)}{dE_i}=\frac{1}{E_\phi}\,, \tag{15}$$

with the total rest frame decay width of ϕ

$$\Gamma_\phi=\frac{g^2}{8\pi}m_\phi\sum_{i,j=1}^3|U_{si}^*U_{sj}|^2\,. \tag{16}$$

The kinematic constraint on the daughter neutrino energies is $E_i, E_j\in[0,m_\phi]$.

Solution of the equations of motion

If we neglect matter effects, the equations of motion (4) can be solved analytically to obtain the electron neutrino flux $\phi_e(L,E)$ appearing in a muon neutrino beam of energy E after a distance L due to oscillations and decay. Neglecting the small mass splittings between the three light neutrino mass eigenstates, $\phi_e(L,E)$ is given by

$$\phi_e(L,E)=\phi_\mu(0,E)\,|U_{e4}|^2|U_{\mu4}|^2\left[1+e^{-\frac{m_4\Gamma_4L}{E}}-2e^{-\frac{m_4\Gamma_4L}{2E}}\cos\left(\frac{\Delta m_{41}^2L}{2E}\right)\right]+|U_{\mu4}|^2\frac{|\langle\nu_e|\nu_F\rangle|^2}{|\langle\nu_F|\nu_F\rangle|^2}\mathcal{I}\,. \tag{17}$$

Here $\phi_\mu(0,E)$ is the initial ν_μ flux, $|\nu_F\rangle=\sum_{i=1}^3U_{si}^*|\nu_i\rangle$ is the superposition of mass eigenstates into which ν_4 decays, and the decay integral $\mathcal{I}$ is

$$\begin{aligned}\mathcal{I}=&\int_{E/(1-x_{\phi4}^2)}^\infty dE_4\left(1-e^{-\frac{m_4\Gamma_4L}{E_4}}\right)\phi_\mu(0,E_4)\sum_j\frac{1}{\frac{m_4}{E_4}\Gamma_4}\frac{d\Gamma^{\rm lab}(\nu_4\to\nu_j\phi)}{dE}\\&+\int_E^\infty dE_\phi\int_{E_\phi}^{E_\phi/x_{\phi4}^2}dE_4\frac{1}{\frac{m_4\Gamma_4L}{E_4}-\frac{m_\phi\Gamma_\phi L}{E_\phi}}\left[\left(1-e^{-\frac{m_\phi\Gamma_\phi L}{E_\phi}}\right)\frac{m_4\Gamma_4L}{E_4}-\left(1-e^{-\frac{m_4\Gamma_4L}{E_4}}\right)\frac{m_\phi\Gamma_\phi L}{E_\phi}\right]\\&\times\frac{1}{\frac{m_4}{E_4}\Gamma_4}\sum_j\left[\phi_\mu(0,E_4)\frac{d\Gamma^{\rm lab}(\nu_4\to\nu_j\phi)}{dE}+\bar\phi_\mu(0,E_4)\frac{d\Gamma^{\rm lab}(\bar\nu_4\to\bar\nu_j\phi)}{dE}\right]\sum_{i,j}\frac{1}{\frac{m_\phi}{E_\phi}\Gamma_\phi}\frac{d\Gamma^{\rm lab}(\phi\to\nu_i\bar\nu_j)}{dE}\,.\end{aligned} \tag{18}$$

Here, $\phi_\mu(0,E)$ and $\bar\phi_\mu(0,E)$ are the initial ν_μ and $\bar\nu_\mu$ fluxes, respectively. A completely analogous equation describes $\bar\nu_e$ appearance.

The physical interpretation of eq. (17) is straightforward: the first term on the right-hand side describes $\nu_\mu\to\nu_e$ oscillations, altered by the removal of neutrinos at energy E due to ν_4 decay. In fact, this contribution matches the result of ref. [64], on invisible ν_4 decay. The second term gives the contribution from neutrinos generated in ν_4 and ϕ decays. The factor $|U_{\mu4}|^2$ arises because ν_4 is the only mass eigenstate that decays. It describes the amount of ν_4 in the ν_μ beam. The factor $|\langle\nu_e|\nu_F\rangle|^2/|\langle\nu_F|\nu_F\rangle|^2$ is the probability of the decay product to be detected as an electron neutrino, and the integral $\mathcal{I}$ controls the energy distribution of the decay products.

Appendix C: Detailed Investigation of the Parameter Space

To supplement fig. 2 and give the reader a broader overview of the preferred parameter regions of decaying sterile neutrinos, we show in figs. 6 and 7 additional slices through the 5-dimensional parameter space.

The color coding in the figure is the same as in fig. 2: the yellow, banana-shaped regions are preferred by MiniBooNE, the large dark red ones by LSND; the orange regions at low $|U_{e4}|^2$ correspond to a global fit to MiniBooNE, OPERA, ICARUS, E776, KARMEN, nuclear beta spectra, and cosmological free-streaming constraints; bright red regions show instead a global fit to MiniBooNE, LSND, OPERA, ICARUS, E776, KARMEN, and nuclear beta spectra, but excluding the free-streaming constraint. Solid lines indicate constraints from OPERA (blue), ICARUS (purple), KARMEN (cyan), E776 (green), nuclear beta decay spectra (black dashed), and free-streaming in the early Universe (black dotted). The region to the right of the lines is excluded.

We observe that, at smaller values of $m_4\Gamma_4$, the allowed parameter regions from short-baseline oscillations (MiniBooNE, LSND, KARMEN) shift towards larger values of $|U_{e4}|^2$ and $|U_{\mu 4}|^2$. In this case, only a small fraction of neutrinos decays before reaching the detector, making the phenomenology more similar to that of $3+1$ models without decay. Strong constraints from beta decay spectra and from cosmology imply that a good global fit cannot be achieved at $m_4\Gamma_4 \ll 1\,\mathrm{eV}^2$.

Regarding the dependence of the fit on m_4, we note that smaller values of m_4 are favored by beta decay spectra, but disfavored by cosmology, in agreement with fig. 3. Exclusion limits from oscillation experiments do not depend on m_4 for $m_4 \gg \mathrm{eV}$.

Comparing fig. 6 with $m_\phi/m_4 = 0.5$ and fig. 6 with $m_\phi/m_4 = 0.9$, we see that it becomes in general more difficult to fit all experiments at smaller m_ϕ/m_4. The reason is that, at small m_ϕ/m_4, the active neutrinos produced in ν_4 and ϕ decays have a harder spectrum. This in particular makes it more difficult to explain the MiniBooNE low-energy excess. In fact, for even smaller values of m_ϕ/m_4, and in particular for nearly massless ϕ (as considered in refs. [21, 22]), the MiniBooNE-preferred region would disappear completely from the plots.

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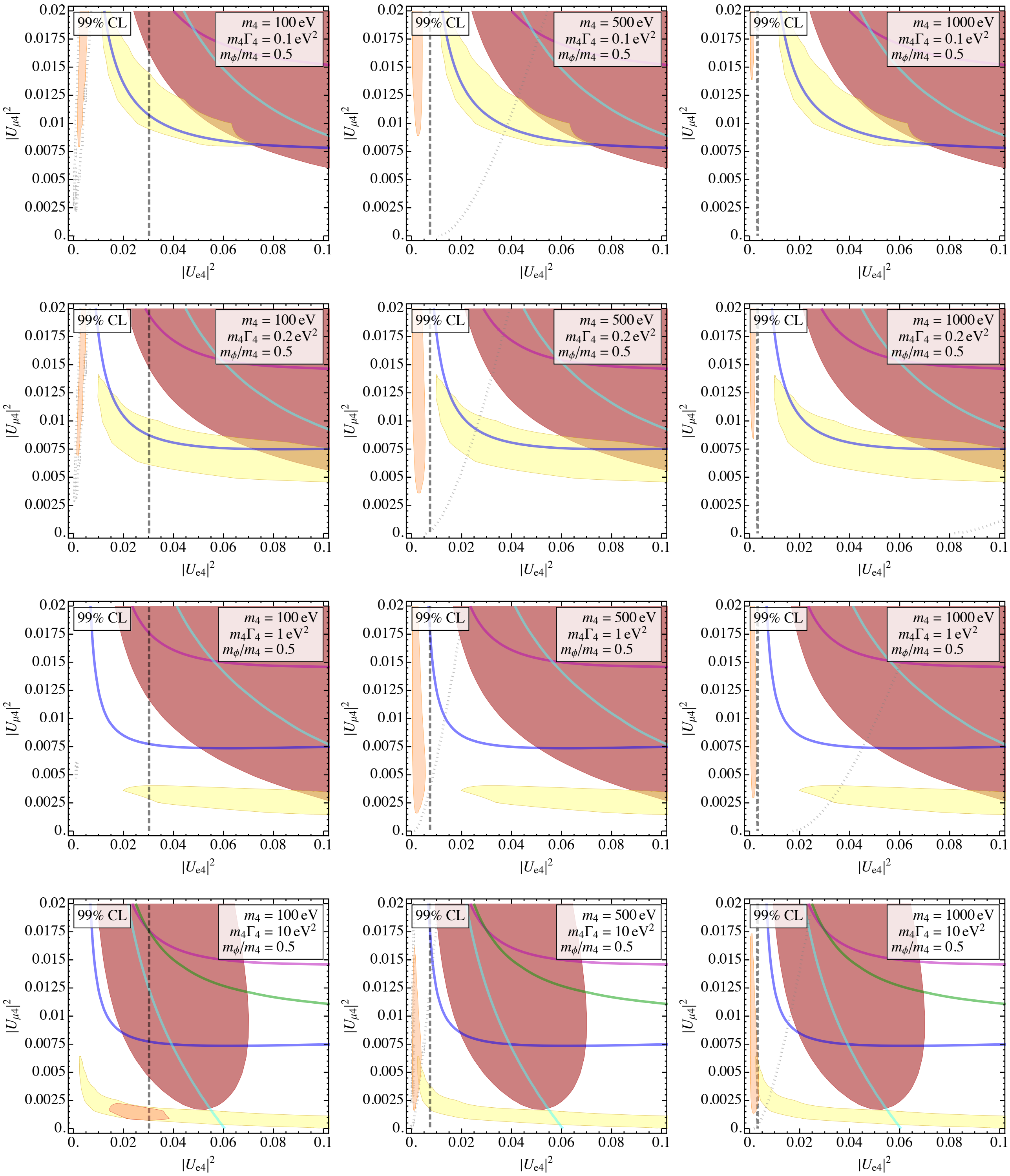

FIG. 6. Slices through the 5-dimensional parameter space of decaying sterile neutrinos at $m_\phi/m_4 = 0.5$ fixed. The color code is the same as in fig. 2.

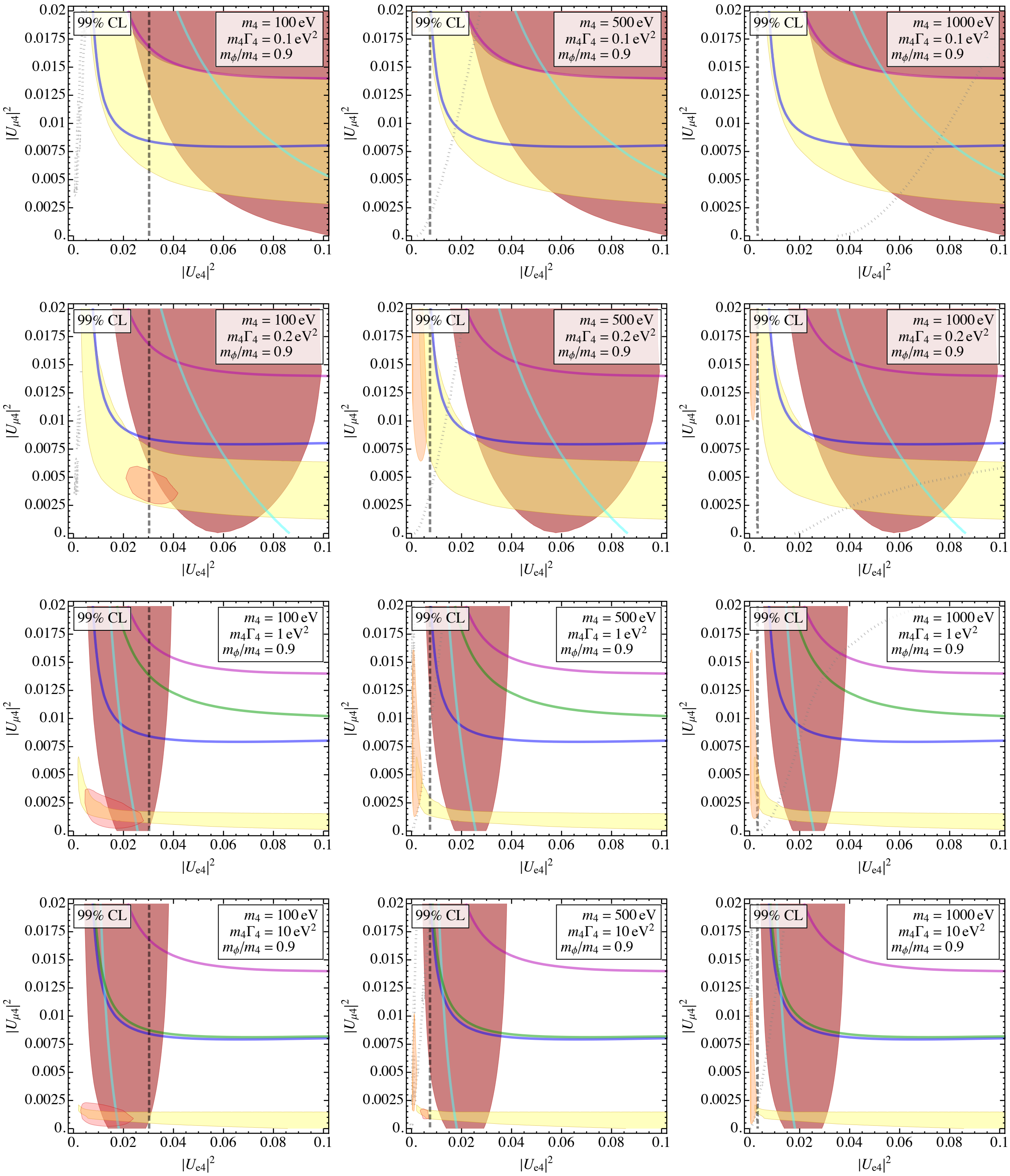

FIG. 7. Slices through the 5-dimensional parameter space of decaying sterile neutrinos at $m_\phi/m_4 = 0.9$ fixed. The color code is the same as in fig. 2.