

Signatures of quasi-Dirac neutrinos in diffuse high-energy astrophysical neutrino data

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Although the sources of astrophysical neutrinos are still unknown, they are believed to be produced by a population of sources in the distant Universe. Measurements of the diffuse, all-sky astrophysical flux can thus be sensitive to flavor- and energy-dependent propagation effects, such as very long-baseline oscillations. These oscillations are present in certain neutrino mass models, such as when neutrinos are quasi-Dirac. Assuming generic models for the source flux, we find that these oscillations can still be resolved even when integrated over wide distributions in source redshift. We use two sets of IceCube all-sky flux measurements, made with muon and all-flavor neutrino samples, to set constraints at the 3σ level on quasi-Dirac mass splittings between $(5 \times 10^{-19}, 8 \times 10^{-19})$ eV². Additionally, we find that when the mass-square splittings are different between generations, the quasi-Dirac scenario can explain tensions between the all-flavor and muon-flavor measurements. We consider systematic uncertainties on the source population and find that our results are robust under alternate spectral hypotheses or physical redshift distributions. Our analysis shows that spectral features in the all-sky neutrino measurements provide strong constraints on massive neutrino scenarios and are sensitive to uncharted parameter space.

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I. INTRODUCTION

The fundamental origin of neutrino oscillations and the precise nature of neutrino mass remain unknown. Currently, the best bet we have to resolve the Dirac versus Majorana nature of neutrinos is by observing the neutrinoless double beta decay ($0\nu\beta\beta$) process [1]. However, there is no guarantee that even next-generation ton-scale $0\nu\beta\beta$ experiments will yield a positive signal if neutrino masses follow the normal ordering, a scenario that is mildly

favoured by global oscillation fits [2]. On the other hand, traditional oscillation-based neutrino experiments are only sensitive to the squared-mass differences between generations, and do not alter chirality, hence, they cannot resolve the distinction between Dirac and Majorana neutrinos. As we will show below, however, measurements of astrophysical neutrino spectra are sensitive to the oscillations of high-energy neutrinos produced by extremely small squared-mass differences, and can thus test a whole class of neutrino mass models with tiny lepton number violation. In fact, this could be the first testable prediction of a class of string theory landscape constructions.

If right-handed neutrino fields ν_R exist, neutrinos may acquire their mass in the same way the charged fermions do, e.g., via a Dirac mass term originating from a Yukawa coupling to the Higgs field. However, since these new fields would be singlets under the standard model (SM) gauge group, no symmetry would then forbid a Majorana mass term for the right-handed fields. Since, in a bottom-up phenomenological approach, the Majorana mass m_R is completely unconstrained, it is theoretically allowed to be

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much smaller than the original Dirac mass term [3–7]. This scenario produces what are called quasi-Dirac neutrinos, and is the focus of this work. Though we follow a phenomenological bottom-up approach, quasi-Dirac neutrinos are also motivated by quantum gravity theories. For example, the Dirac nature of neutrinos has been predicted by string theory landscape constructions [8–12] and sources of lepton number violation, such as small m_R terms, expected from quantum gravity theories. This makes the quasi-Dirac regime one of the few testable predictions of string theory.

The only experimental way to directly probe quasi-Dirac models is by searching for oscillations in the spectra of neutrinos from astrophysical sources, such as the Sun [13–19], supernovae [20,21], high-energy cosmic sources [22–34], or relic neutrinos [35]. These oscillations are driven by hyperfine mass-squared differences, $\delta m_k^2 \propto m_R$. Stringent upper limits on $\delta m_{1,2}^2 \lesssim 10^{-12}$ eV² have been derived using solar neutrinos [15,18]. These solar limits supersede a previously derived constraint from big bang nucleosynthesis (BBN), $\delta m_k^2 \lesssim 10^{-8}$ eV² [36,37]. Both these limits are derived assuming maximal active-sterile neutrino mixing in the quasi-Dirac scenario. If the mixing is nonmaximal, the solar δm^2 limits can be much weaker [18,38]. Moreover, the solar neutrino data are not sensitive to δm_3^2 due to the small electron component in ν_3 , and the limits from atmospheric, accelerator, and reactor neutrino data are very weak, $\delta m_3^2 \lesssim 10^{-5}$ eV² [16], due to the much shorter baselines.

The identification of a few point sources of astrophysical neutrinos [39] provides sensitivity to quasi-Dirac neutrinos with $\delta m_k^2 \in [10^{-21}, 10^{-16}]$ eV² [22], see also Refs. [40,41]. However, the current prospective for exploring the quasi-Dirac parameter space in this mass range using point sources is limited by the small number of identified sources, poor resolution on the muon neutrino energy, and lack of flavor information. These problems do not allow high-significance statements about quasi-Dirac neutrinos using point sources at present.

In this work, for the first time, we use the diffuse astrophysical neutrino flux measured by IceCube in multiple channels and flavors to explore quasi-Dirac models with mass-squared differences $\delta m_k^2 \in [10^{-21}, 10^{-16}]$ eV². This strategy has two advantages over point source-based analyses: firstly, measurements in multiple channels with different flavor combinations are available, and secondly, the astrophysical neutrino sample size is much larger. Furthermore, as we will show, for physically motivated choices of the spatial distribution of astrophysical sources, the quasi-Dirac oscillation probability is not fully smeared out. It is, therefore, possible to search for oscillation-induced disappearance dips in the astrophysical neutrino spectrum.

We study two classes of quasi-Dirac scenarios: a simpler one-parameter model without flavor dependence, and a more complex two-parameter model that introduces flavor-dependent oscillation effects. As shown in Fig. 1, our results in the

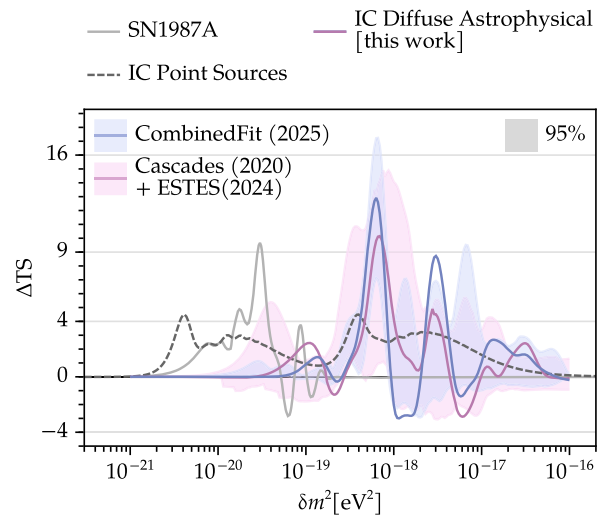


FIG. 1. Constraints on the quasi-Dirac parameter space. The test statistic difference with respect to the broken power-law null hypothesis, as a function of the quasi-Dirac mass-squared difference δm^2 , is plotted for analyses based on IceCube’s 2025 CombinedFit [49] (blue) and on a combination of Cascades (2020) [50] and ESTES (2024) [43] results (purple). Brazil bands indicate the regions contained by 95% of realizations drawn from the best-fit broken power law (BPL) spectrum in [49]. Previous constraints from Ref. [20] and sensitivities from Ref. [22] are shown in gray.

first class considered find that measurements of the diffuse astrophysical flux constrain regions of quasi-Dirac squared-mass difference δm^2 from $(5\text{--}8) \times 10^{-19}$ eV² at the 3σ global significance for the first time. These constraints are driven by the absence of disappearance signatures in the spectrum around a few tens of TeV. The second scenario we consider is interesting because flavor-dependent disappearance effects could help explain tensions in the astrophysical spectrum measured by different neutrino flavor combinations. In particular, all-flavor measurements of the astrophysical flux now find strong preference for a break in the spectrum at 30 TeV [42], which is not present in muon neutrino-only measurements [43]. In our analysis, we find that flavor-dependent quasi-Dirac effects can jointly explain all-flavor and muon-flavor measurements without needing to introduce a spectral break. This implies that quasi-Dirac neutrino mass models impact the interpretation of astrophysical neutrino flux measurements.¹

II. THEORY OF QUASI-DIRAC NEUTRINOS

It is unknown whether the neutrino mass term is Majorana, Dirac, or a mixture of both [3,4]. In particular, if the Majorana mass term is much smaller than the Dirac mass, neutrinos are quasi-Dirac [5–7], which are

¹For other explanations of the track vs cascade tension in the previous IceCube data, see, e.g., Refs. [44–48].

fundamentally Majorana fermions but behave like Dirac in most experimental settings. In the quasi-Dirac scenario, the lepton number is only slightly violated, and lepton number-violating observables such as $0\nu\beta\beta$ are strongly suppressed. Ordinary neutrinos have three generations of mass states, ν_i , whose decomposition into flavor states, ν_α , is described by the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrix. If neutrinos are quasi-Dirac, each generation has a further hyperfine mass splitting, δm_k^2 , which produces two mass states per generation, $\nu_k^\pm$, which are orthogonal equal combinations of active (left-handed) and sterile (right-handed) states. In the same way that ordinary neutrinos undergo flavor oscillations after traveling distances $L \sim E/\Delta m^2$, quasi-Dirac neutrinos undergo active-sterile oscillations over much longer distances $L \sim E/\delta m^2$, where E is the neutrino energy. These ultralong-baseline oscillations are the only distinguishing signature of quasi-Dirac neutrinos.

The theoretical and model-building aspects of quasi-Dirac neutrinos have been extensively discussed in the literature, see, e.g., Refs. [15,22,29,51–59]. It is interesting to note that certain string theory landscape constructions, such as Swampland, predict that neutrinos are necessarily Dirac-like particles [8–12]. Additionally, we expect that all global symmetries including the lepton number are ultimately broken by quantum gravity, turning Dirac neutrinos into quasi-Dirac. In fact, any model where neutrinos start as Dirac particles with conserved lepton number could receive nonrenormalizable quantum gravity corrections, which generate small δm^2 via higher-dimensional lepton number-violating operators suppressed by the Planck scale, thus making neutrinos naturally quasi-Dirac. The lowest lepton number-violating dimension-5 Weinberg operator typically gives $\delta m \sim v^2/M_{\text{Pl}} \sim 10^{-5}$ eV, which implies $\delta m^2 \sim 2m\delta m \sim 10^{-6}$ eV², where $m \sim 0.05$ eV is the light neutrino mass scale, v is the electroweak scale, and M_{Pl} is the Planck scale. However, in models with a gauged lepton number, where the dimension-5 operator is forbidden and lepton number breaking is induced only at dimension-6 or higher, δm^2 is further suppressed by higher powers of M_{Pl} , thus allowing for a wider range of δm^2 values consistent with the existing solar neutrino constraints. For instance, in the gauged $B-L$ model considered in Ref. [22], $\delta m \sim v^2 v_{BL}/M_{\text{Pl}}^2$, where v_{BL} is the $B-L$ breaking scale, and depending on its value, δm^2 can be in the range of $(10^{-22} \text{--} 10^{-12})$ eV².

Quasi-Dirac neutrinos could also help explain some features of astrophysical and cosmological observations. For example, small δm^2 values could also be linked to the observed baryon asymmetry of the Universe [58,60]. Quasi-Dirac neutrinos have also been suggested as solutions to the excess in the diffuse radio background [61,62]. Thus, from both theoretical and phenomenological perspectives, the search for quasi-Dirac neutrinos is highly motivated.

III. QUASI-DIRAC OSCILLATIONS ON ASTROPHYSICAL SCALES

A quasi-Dirac neutrino mass model modifies standard flavor oscillations by introducing additional frequencies beyond those associated with the “solar” ($\Delta m_{21}^2 \approx 7.5 \times 10^{-5}$ eV²) and “atmospheric” ($\Delta m_{31}^2 \approx 2.5 \times 10^{-3}$ eV²) mass-squared differences. These new oscillation components convert mostly active neutrino mass eigenstates, ν_k , into their sterile counterparts with frequencies proportional to δm_k^2 . For any combination of neutrino energy, E , and propagation distance, L , suitable for probing allowed quasi-Dirac mass-squared differences, δm^2 , the standard oscillation terms containing Δm^2 will average out due to limited experimental resolution. In this regime, the oscillation probability is given by [22]

$$P_{\alpha\beta} = \frac{1}{2} \sum_{k=1}^3 |U_{\beta k}^* U_{\alpha k}|^2 \left[1 + \cos\left(\frac{\delta m_k^2 L_{\text{eff}}}{2E}\right) \right]. \quad (1)$$

Here, E is the neutrino energy as measured on Earth, and

$$L_{\text{eff}} = \int \frac{dz}{H(z)(1+z)^2} \quad (2)$$

is the effective distance traversed, which accounts for neutrino propagation in an expanding Universe and depends on the Hubble expansion rate, $H(z)$.²

The diffuse neutrino flux is produced by a population of sources distributed across cosmological time and space. Without specifying the exact nature of these sources, we assume that they are distributed in redshift according to some evolution function $\rho(z)$, e.g., proportional to the star formation rate density (SFRD). In this scenario, the number of neutrino sources within a redshift interval $[z, z+dz]$ is given by [64]

$$dN(z) = \rho(z) \times 4\pi D^2 \frac{dD}{dz} dz, \quad (3)$$

where $D(z) = \int dz/H(z)$ is the comoving distance.

²Note that, unlike standard oscillatory terms, quasi-Dirac oscillations are not necessarily washed out by extragalactic baselines. This is because the coherence length is

$$L_{\text{coh}} = \frac{4\sqrt{2}E^2}{|\delta m_k^2|} \sigma_x \approx 18 \text{ Gpc} \left(\frac{E}{10 \text{ TeV}} \right)^2 \left(\frac{10^{-19} \text{ eV}^2}{\delta m_k^2} \right) \left(\frac{\sigma_x}{10^{-19} \text{ m}} \right).$$

Therefore, for benchmark values of 10 TeV and $\delta m_k^2 = 10^{-19}$ eV², the coherence length is comparable to the radius of the observable Universe, even for the size of the wave packet $\sigma_x \sim 10^{-19}$ m, orders of magnitude smaller than the smallest wave packets ever considered [63].

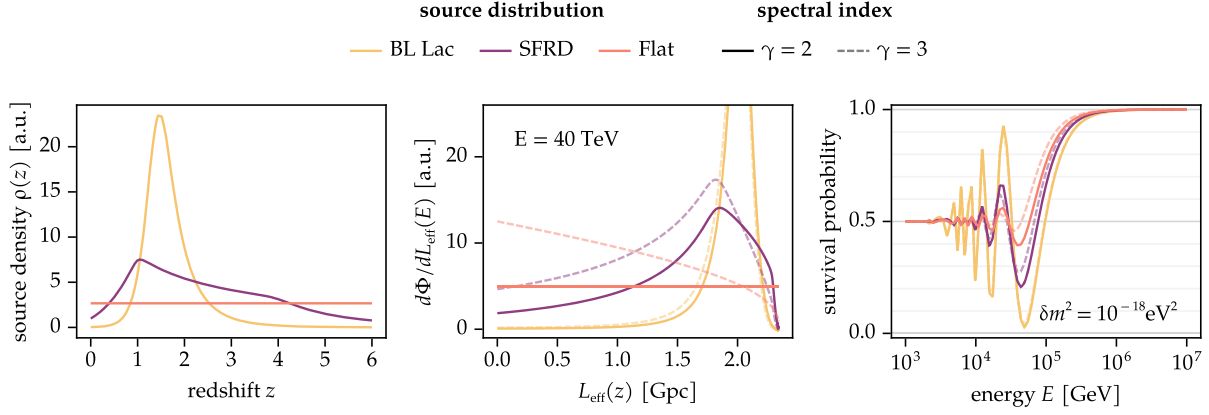


FIG. 2. Source evolution and integrated oscillation probability. Left: three choices for the source evolution in redshift: a flat distribution, the SFRD of Ref. [65], and the distribution of BL Lacs from Ref. [66]. Middle: distribution of neutrinos with energy $E = 40$ TeV over the effective distance (L_{eff}), assuming a given redshift evolution function $\rho(z)$ and power-law emission spectrum with index γ . Right: quasi-Dirac survival probability as a function of detected neutrino energy, integrated over the population of sources at all redshifts.

In this work, we assume that sources at all redshifts emit neutrinos with the same spectrum and in the same ratio of flavors. In this case, the flux of neutrinos of flavor α at Earth from a source at redshift z is given by

$$\Phi_\alpha(E, z) = f_\alpha \phi_0(E(1+z)) \times \frac{(1+z)^2}{4\pi D_L^2}. \quad (4)$$

In this expression, f_α is the fraction of neutrinos produced in flavor α , and ϕ_0 is the common functional form for the emitted neutrino spectra.

Including quasi-Dirac oscillations, the contribution to the total diffuse flux from a given redshift interval is

$$d\Phi_\beta(z, E) = \sum_\alpha P_{\alpha\beta}(E, z) \Phi_\alpha(E, z) dN(z). \quad (5)$$

Thus, the total flux at Earth as a function of energy is, therefore, given by the integral over redshift,

$$\Phi_\beta(E) = \int dz \sum_\alpha P_{\alpha\beta}(E, z) \times f_\alpha \phi_0(E(1+z)) \times \frac{\rho(z)}{H(z)}. \quad (6)$$

In Fig. 2, we show a range of possibilities for the source redshift evolution function $\rho(z)$: a flat distribution in z , the star formation rate density function from Ref. [65], which grows for small z as $(1+z)^3$, and the distribution of BL Lacs from Ref. [66], which grows for small z as $(1+z)^5$. These distributions all correspond to mean effective distances L_{eff} on the order of Gpc, and thus produce first oscillation minima at tens of TeV for a mass-squared difference $\delta m^2 = 10^{-18} \text{ eV}^2$. The SFRD and BL Lac distributions, which grow steeply with z , produce a much more pronounced disappearance effect than the flat distribution. The oscillation curve also changes with different

choices of the source emission spectrum ϕ . A softer flux, with a power-law spectral index $\gamma = 3$ instead of 2, produces a smaller disappearance signature, since more of the neutrinos at a given energy E will be coming from nearby sources.

IV. ANALYSIS

Quasi-Dirac oscillations can modify both the energy spectrum and the flavor composition of propagated neutrinos. If the mass-squared differences δm_k^2 are different for each mass state pair k , the different flavor states α will be affected by different combinations of oscillations, introducing new flavor dependence to the neutrino spectrum. Alternatively, if the mass-squared differences are all the same $\delta m_k^2 = \delta m^2$, there will be no flavor-dependent effects, but the magnitude of the oscillation effects will be maximized for all flavors.

To be sensitive to both energy-only and flavor-dependent effects, we combine IceCube results from two statistically independent datasets, *Cascades* and *ESTES*, that select events with different morphologies. The *Cascades* sample [50] is a selection of cascade events collected over 6.5 years of data taking. The *ESTES* sample [43] is a selection of track events whose production vertex is within the detector volume (“starting tracks”) collected over 10.3 years of data taking. Since tracks in IceCube are produced primarily by muon neutrinos, with a subdominant contribution from tau neutrinos that decay to muons, whereas cascades are produced by all flavors in comparable fractions, these two selections are sensitive to different flavor combinations.

Additionally, IceCube has recently published its own CombinedFit analysis of the astrophysical flux, which combines an updated version of the *Cascades* sample (11.5 years) with a northern sky track selection [42,49].

Due to the increased lifetime and better control of systematics, this new result is much more sensitive to the shape of the astrophysical flux. However, because the result is reported jointly for the cascade and track event selections, it cannot be used to study flavor-dependent effects. We, therefore, study only the equi-splitting $\delta m_k^2 = \delta m^2$ scenario with the CombinedFit results.

IceCube's analyses model the astrophysical neutrino spectrum by splitting the spectrum piecewise into a series of energy bins, each described by an E^{-2} power law with independent normalization. The analyses then report each segment normalization and their $\pm 1\sigma$ variations. These piecewise fluxes can be converted back into the most likely number of astrophysical neutrinos per bin by taking the product with the detector acceptance,

$$N_i^S = T^S \int_{E_i}^{E_{i+1}} dE \left[E^{-2} \Phi_i^S \sum_{\alpha} A_{\alpha}^S(E) \right]. \quad (7)$$

In this expression, the index i indicates the energy bin, the index S indicates the sample, T^S is the total lifetime of the sample, Φ_i^S is the reported per-bin normalization, and A_{α}^S is the flavor-dependent, zenith-averaged effective area of the sample. Similarly, we can calculate approximate $\pm 1\sigma$ variations on the number of events N_i^S by adjusting the per-bin normalization Φ_i^S . These approximate observed event counts can then be compared to the predictions of an astrophysical flux model with quasi-Dirac oscillations,

$$\mu_i^S(\delta m_k^2, \vec{\eta}) = T^S \int_{E_i}^{E_{i+1}} dE \left[\sum_{\alpha} \Phi_{\alpha}(E | \delta m_k^2, \vec{\eta}) A_{\alpha}^S(E) \right]. \quad (8)$$

The flux model $\Phi_{\alpha}(E)$ depends on the quasi-Dirac parameters, δm_k^2 , as well as on the assumed source redshift evolution function $\rho(z)$, the source flavor ratio f_{α} , and the spectrum at the source, ϕ , which itself depends on additional parameters, e.g., a spectral index. We collectively denote these as the flux parameters $\vec{\eta}$, and marginalize over them in our analysis, as described in Appendix A. When calculating flavor-dependent oscillation probabilities using Eq. (1), we use the PMNS matrix values reported in Ref. [67]. Moreover, we use the results of Ref. [68] for the parameters of our Λ CDM cosmology.

To determine preferred and disfavored regions of δm^2 parameter space, relative to the nonquasi-Dirac hypothesis, we use the following test statistic (TS):

$$\text{TS}(\delta m_k^2) = \min_{\vec{\eta}} \sum_S \sum_i -2 \log \mathcal{L}_i(\mu_i^S(\delta m_k^2, \vec{\eta})), \quad (9)$$

where $\mathcal{L}_i$ is an approximation of the IceCube likelihood function. We construct this approximation differently in two regimes: for lower-energy bins with $E_i < 100$ TeV, we use a Gaussian likelihood with mean N_i^S and standard deviation σ_i^S , while for higher-energy bins above 100 TeV we use a

calibrated Poisson likelihood, which is scaled so that $-2 \log \mathcal{L}_i(N_i^S) = 0$ and $-2 \log \mathcal{L}_i(N_i^S \pm \sigma_i^S) = 1$. This likelihood approximation yields good agreement with IceCube results: it accurately captures the low-energy regime, where the data are dominated by backgrounds and systematics, and remains reliable at high energies, where statistical uncertainties prevail. More details on this likelihood function, as well as a comparison of published IceCube fits and the results using our likelihood, can be found in Appendix B. Assuming Wilks' theorem, we construct our confidence regions by assuming that the difference $\Delta \text{TS}(\delta m_k^2) = \text{TS}(\delta m_k^2) - \text{TS}(0)$, with the appropriate sign for preference or exclusion, follows a χ^2 distribution.

V. RESULTS: EQUAL QUASI-DIRAC SQUARED-MASS DIFFERENCES

We begin with our results for the energy-only analysis, in which we assume a single mass-squared difference, δm^2 , for all three generations. In this scenario, the initial flavor ratio has no effect on the quasi-Dirac signal, and we, therefore, assume a pion-decay production mechanism. For our main result, shown in Fig. 1, we assume a broken power-law flux hypothesis for ϕ , which is strongly preferred over a single power law in the most recent IceCube astrophysical flux results [42,49], and redshift evolution function $\rho(z)$ given by the star formation rate density model of [65]. In this case, we find that IceCube's CombinedFit results disfavor $\delta m^2 \in [5.0 - 7.5] \times 10^{-19} \text{ eV}^2$ at the 3σ level, and our combination of the Cascades and ESTES results disfavor a similar region $\delta m^2 \in [5.9 - 7.9] \times 10^{-19} \text{ eV}^2$ at the same level. These statements reflect the fact that the piecewise fluxes reported by IceCube are well described by an unbroken power law in the energy range just below the break at 50 TeV, as can be seen from the lower middle panels in Figs. 4(a) and 4(b). The constraints we obtain are stronger than the projected sensitivities of individual IceCube point sources [22], shown in Fig. 1 by the dark gray dashed curve, and are slightly stronger than those obtained from SN1987A neutrino data [20]. We do not find any significant preference for a quasi-Dirac hypothesis.

To characterize the significance of our results, we generated “Brazil bands” by generating 100 pseudoexperiments assuming the best-fit BPL flux reported in IceCube's CombinedFit results [42]. We find that the disfavored region we get from the published results is contained in the 95% band. Elsewhere, excursions from the band correspond to individual bins that deviate from the BPL expectation. Thus, the result of our analysis is within expectations of statistical fluctuations. The ΔTS curve based on the Cascades and ESTES data shows a slight minimum around $6 \times 10^{-18} \text{ eV}^2$, which is not contained within the 95% band. This is due to the deficit of events in the two bins in the Cascades results just below 1 PeV. A similar

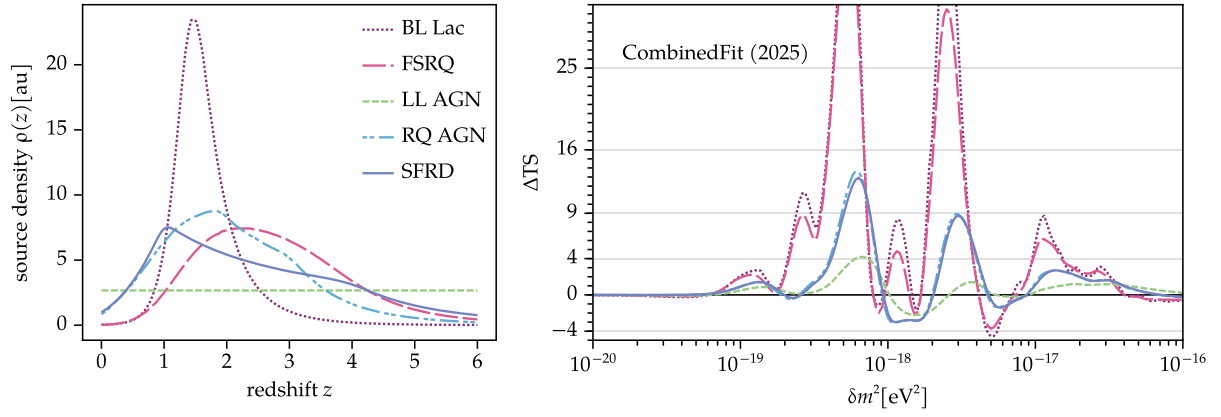


FIG. 3. Constraints on the quasi-Dirac mass-squared difference, under different choices for the source redshift evolution $\rho(z)$. Left: redshift distributions of four source populations from Ref. [66], as well as the SFRD of Ref. [65]. Note that as discussed in the text, models in which neutrino sources are distributed uniformly over redshift, such as that of the LL AGNs in this plot, are disfavored by source studies [64]. Right: the resulting TS curves calculated based on the CombinedFit flux points, using a BPL emission spectrum.

but less significant effect can also be seen in the CombinedFit curve. The oscillatory patterns in the ΔTS curve are artifacts of the binned representation of the astrophysical fluxes and can be improved in internal experimental analyses.

When we change the choice of redshift evolution function $\rho(z)$, we find that the strength at which we can disfavor portions of quasi-Dirac parameter space can change, but the regions where significant conclusions are drawn remain substantially the same, as can be seen in Fig. 3. For this study, we considered four physical redshift evolution models from [66], representing four classes of neutrino sources: BL Lacs, flat spectrum radio quasars (FSRQs), low-luminosity (LL), active galactic nucleus (AGNs), and radio-quiet (RQ) AGNs. These models also represent a range of possible redshift scalings: the BL Lac and FSRQ models both scale as $(1+z)^5$ for small z , whereas the RQ AGN model scales, like the SFRD, as $(1+z)^3$, and the LL AGN model is flat in z . This is important because, as can be seen in Fig. 2, the scaling with z determines the source distribution over the quasi-Dirac oscillation length, L_{eff} , and thus the extent to which the quasi-Dirac disappearance signal is resolved or smeared out.

We find that the models with similar scaling behavior result in very similar TS curves, and thus similar excluded

regions (see Table I). The BL Lac and FSRQ models, which grow as $(1+z)^5$, both very strongly exclude regions around $(3.5-7.2) \times 10^{-19}$ and $(1.8-3.5) \times 10^{-18} \text{ eV}^2$, due to nonobservation of disappearance features at around 50 TeV or 100 TeV. Conversely, the RQ AGN and SFRD models, which both grow as $(1+z)^3$, predict less pronounced disappearance features and only disfavor $\delta m^2 \in (5-7.5) \times 10^{-19} \text{ eV}^2$ at the 3σ level. Finally, the LL AGN model, which is flat in z , produces a very limited disappearance signal, and thus cannot disfavor any δm^2 region at 3σ . However, we note that such flat distributions appear to be disfavored by external analyses of IceCube's high-energy neutrino data, since they would predict more bright sources than have been observed [64]. Thus, when restricting ourselves to source distributions preferred by neutrino source searches, the astrophysical spectrum remains sensitive to long-baseline oscillations.

We also considered how our results are affected by different choices for the source emission spectrum ϕ . For this comparison, we consider three analytic functions commonly used to characterize the astrophysical flux: a single power law (SPL), a single power law with an exponential cutoff at high energies (SPE), and a BPL. The three TS curves, as well as 95% Brazil bands generated under the CombinedFit BPL flux, are shown in Fig. 5.

TABLE I. Limits on the squared-mass difference. These results assume the source emission spectrum is a BPL.

Sample(s)	Redshift dist. $\rho(z)$	3σ region(s) $[\text{eV}^2]$
CombinedFit	SFRD [65]	$(5.0-7.5) \times 10^{-19}$
Cascades + ESTES	SFRD [65]	$(5.9-7.9) \times 10^{-19}$
CombinedFit	BL Lac [66]	$(2.4-3.0), (3.5-7.2), (18-35) \times 10^{-19}$
CombinedFit	FSRQ [66]	$(3.5-7.2), (19-35) \times 10^{-19}$
CombinedFit	LL AGN [66]	
CombinedFit	RQ AGN [66]	$(4.9-7.5) \times 10^{-19}$

TABLE II. Best-fit parameters, for each combination of IceCube results used, assuming a single mass-squared difference δm^2 . These results use the SFRD [65] for the source redshift evolution $\rho(z)$.

Sample(s)	Flux model ϕ	δm^2 [eV ²]	TS	d.o.f.	p-value
CombinedFit					
	SPL	<i>Null</i>	35.67	11	0.02%
	SPL	7.2×10^{-20}	18.93	10	4.11%
	SPE	<i>Null</i>	22.82	10	1.14%
	SPE	6.9×10^{-20}	13.91	9	12.57%
	BPL	<i>Null</i>	9.17	9	42.20%
	BPL	1.2×10^{-18}	6.17	8	62.84%
Cascades + ESTES					
	SPL	<i>Null</i>	28.10	16	3.07%
	SPL	1.9×10^{-19}	25.10	15	4.87%
	SPE	<i>Null</i>	26.67	15	3.16%
	SPE	6.3×10^{-18}	22.64	14	6.63%
	BPL	<i>Null</i>	25.06	14	3.39%
	BPL	6.0×10^{-18}	22.18	13	5.26%

Our fit reproduces the $>4\sigma$ preference for BPL over an SPL reported by IceCube’s CombinedFit analysis [42] in the null limit ($\delta m^2 \rightarrow 0$). We find that this gap between the spectral hypotheses cannot be fully alleviated by introducing quasi-Dirac oscillations. The best SPL+quasi-Dirac model fit produces a $TS = 18.93$ for 10 degrees of freedom, which is a considerably worse fit than the BPL result, $TS = 9.17$ for 9 degrees of freedom (see Table II). We also observe that the tension with the SPL hypothesis substantially worsens the quasi-Dirac+SPL fit in the parts of parameter space that are disfavored, above $\simeq 3 \times 10^{-19}$ eV². The TS curves for all three spectral shapes are consistent with the 95% Brazil band produced under the CombinedFit BPL best-fit spectrum.

The BPL fit to the combined Cascades and ESTES flux measurements, shown in Fig. 6, improves over a SPL by only three TS units. This is due to the smaller data sample and larger systematic errors at lower energies compared to the CombinedFit. In this case, we find that the SPL+quasi-Dirac and SPE+quasi-Dirac models, at $\delta m^2 = 1.9 \times 10^{-19}$ eV², are able to describe the reported fluxes equally as well as a BPL without quasi-Dirac effects. Additionally, a mass-squared difference of $\delta m^2 \simeq 6 \times 10^{-18}$ eV² is able to improve the fit in the energy bins near 1 PeV, where there is a deficit in Cascades. However, neither of these improvements are statistically significant. As we found in the CombinedFit study, the additional degrees of freedom in the BPL hypothesis can absorb some of the tension induced by quasi-Dirac mass-squared differences between 3×10^{-19} and 5×10^{-18} eV².

We conclude that the choice of emission spectrum ϕ does not change the region in δm^2 this analysis is sensitive to, but can modify the strength of statements in either direction. For this reason, we chose the BPL model, which

in the null scenario describes the data best, and produces the most conservative limits, for our main result.

VI. RESULTS: MULTIPLE QUASI-DIRAC SQUARED-MASS DIFFERENCES

We search for flavor-dependent quasi-Dirac signals by combining the Cascades and ESTES measurements. To restrict the available parameter space, we consider only scenarios with two distinct squared-mass differences, e.g., $\delta m_1^2 = \delta m_2^2 \neq \delta m_3^2$. For this study, we assume a BPL for the emission spectrum, and use the SFRD for the source distribution over redshift. We consider three possible scenarios for the flavor ratio at the source: pion decay (1, 2, 0), muon damping (0, 1, 0), and neutron dominated (1, 0, 0). In both the pion-decay and neutron-dominated scenarios, oscillations induced by δm_1^2 disproportionately affect cascades, whereas those induced by δm_3^2 disproportionately affect tracks. In the muon-damped scenario, the electron fraction in both tracks and cascades is substantially reduced, so δm_1^2 oscillations have very little impact on either morphology.

As discussed in the introduction, the Cascades and ESTES measurements of the diffuse flux differ at low energies, below 30 TeV, and at high energies, around 1 PeV. These two measurements are plotted in the lower left panel of Fig. 4(b). As can be seen in the figure, at low energies the Cascades spectrum trends downwards, suggesting a spectral break, whereas the ESTES remain more consistent with the trend at higher energies. All flavor measurements from cascade events also drive the overall $>4\sigma$ preference for a spectral break in the newest IceCube analyses of Ref. [42], including the CombinedFit. These flavor-dependent spectral differences are difficult to explain using just astrophysical production mechanisms, but could be

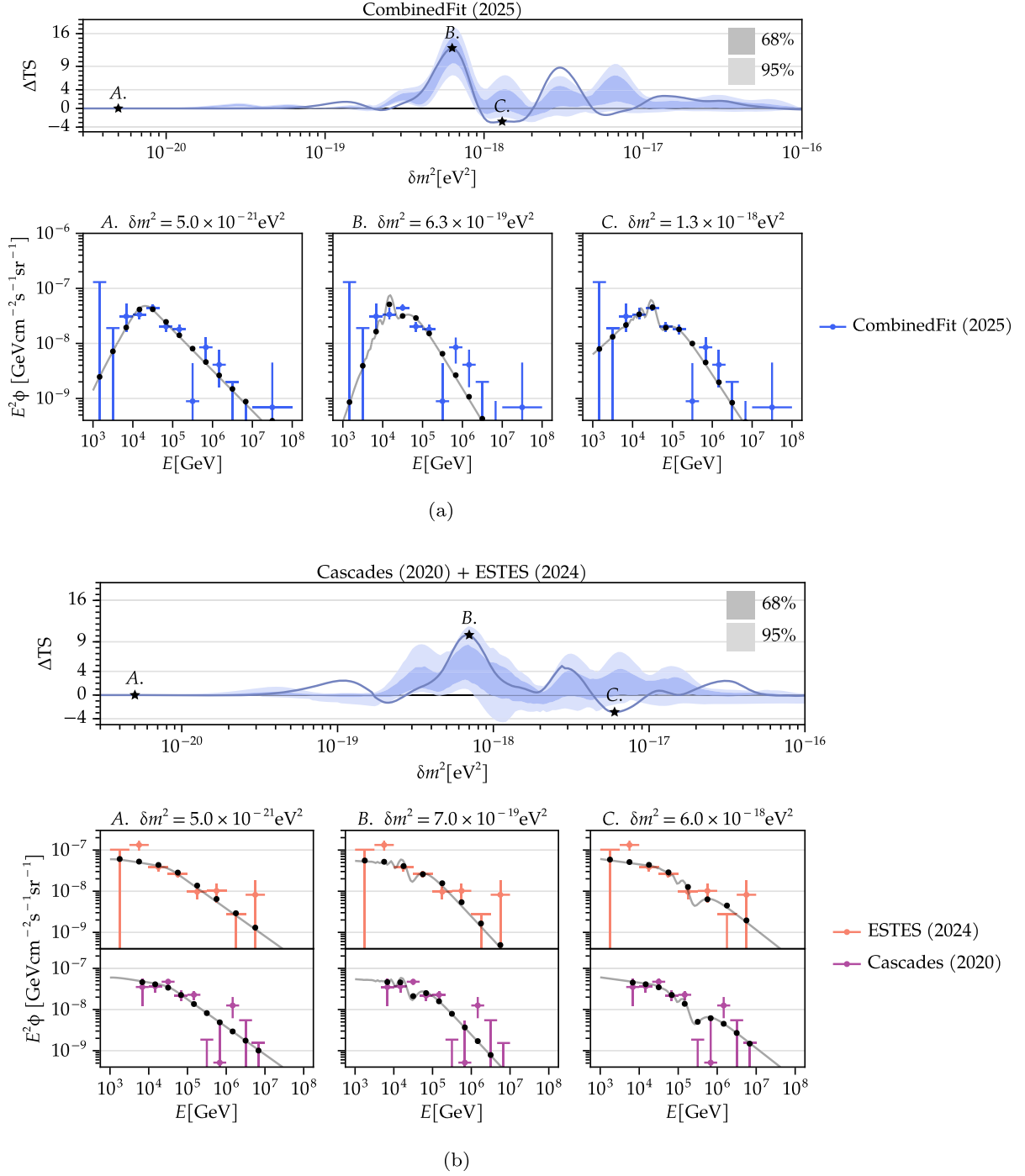


FIG. 4. Quasi-Dirac flux hypotheses, compared with IceCube measurements. Each subfigure is arranged in two parts as follows. Below: the best-fit quasi-Dirac flux hypotheses (gray lines) for three different values of the mass splitting δm^2 , compared with the IceCube measurements. The black points indicate the predicted piecewise flux in each bin according to the quasi-Dirac model. Above: the test statistic difference as a function of δm^2 . The shaded bands indicate the 68% and 95% most probable regions, generated by sampling from the published CombinedFit best-fit BPL flux. (a) Flux predictions and IceCube CombinedFit measurements. (b) Flux predictions and IceCube Cascades + ESTES measurements.

indicative of new propagation physics, such as quasi-Dirac scenarios.

The fit results for three different source flavor assumptions and two different emission hypotheses are listed in

Table III. We find that a quasi-Dirac model with a single power-law emission spectrum can explain the joint measurements better than a broken power law. Assuming a pion-decay initial flavor scenario, the SPL quasi-Dirac

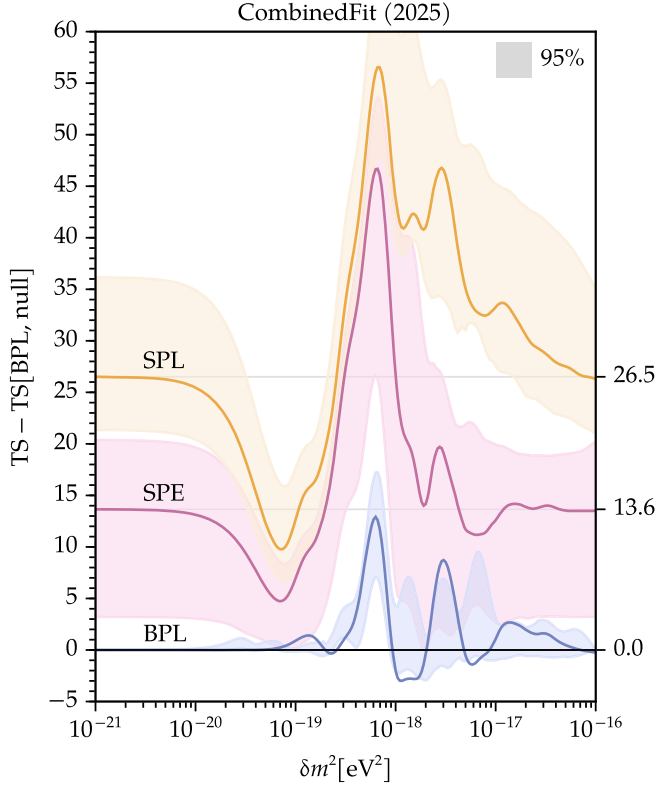


FIG. 5. The test statistic difference between quasi-Dirac models with different source emission models ϕ and the BPL null. The curves shown are calculated based on the CombinedFit data, and the bands indicate the regions contained by 95% of realizations from the published CombinedFit best-fit BPL flux.

model results in a chi-square test statistic of 22.18 units for 14 degrees of freedom, corresponding to a p-value of 7.5%, whereas the BPL model without quasi-Dirac effects produces a test statistic of 25.06 units for 14 degrees

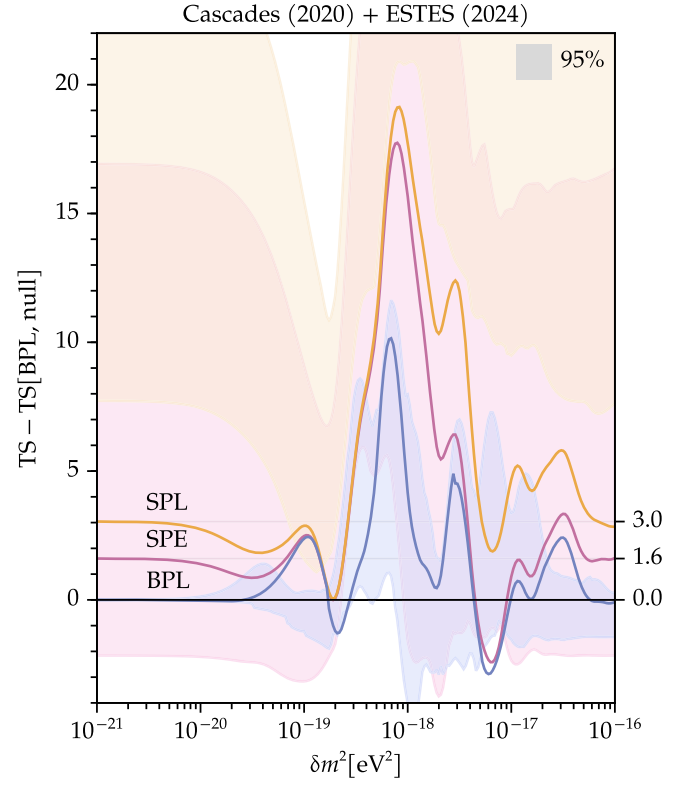


FIG. 6. The test statistic difference between quasi-Dirac models with different source emission models ϕ and the BPL null. The curves shown are calculated based on the combined Cascades and ESTES data, and the bands indicate the regions contained by 95% of realizations from the published CombinedFit best-fit BPL flux.

of freedom, corresponding to a 3.4% p-value. Adding quasi-Dirac effects to a BPL emission spectrum improves the fit further still ($TS = 20.00$) at the cost of 2 additional degrees of freedom. This improvement is illustrated in

TABLE III. Best-fit parameters for fits assuming two distinct mass-squared differences. These results are based on the combined Cascades and ESTES flux points, and use the SFRD [65] for the source redshift evolution function $\rho(z)$.

Flavor ratio	Flux model ϕ	δm_1^2 [eV ²]	δm_2^2 [eV ²]	δm_3^2 [eV ²]	TS	d.o.f.	p-value
Pion decay (1, 2, 0)	SPL	Null	Null	Null	28.10	16	3.07%
Muon-damped (0, 1, 0)	SPL	Null	Null	Null	34.03	16	0.54%
Neutron dom. (1, 0, 0)	SPL	Null	Null	Null	41.58	16	0.05%
Pion decay (1, 2, 0)	SPL	2.0×10^{-19}	$= \delta m_1^2$	5.6×10^{-18}	22.18	14	7.50%
Muon-damped (0, 1, 0)	SPL	1.8×10^{-19}	$= \delta m_1^2$	6.3×10^{-18}	28.30	14	1.30%
Neutron dom. (1, 0, 0)	SPL	2.0×10^{-19}	2.5×10^{-20}	$= \delta m_2^2$	30.80	14	0.59%
Pion decay (1, 2, 0)	BPL	Null	Null	Null	25.06	14	3.39%
Muon-damped (0, 1, 0)	BPL	Null	Null	Null	32.52	14	0.34%
Neutron dom. (1, 0, 0)	BPL	Null	Null	Null	33.35	14	0.26%
Pion decay (1, 2, 0)	BPL	2.0×10^{-19}	5.6×10^{-18}	$= \delta m_2^2$	20.00	12	6.71%
Muon-damped (0, 1, 0)	BPL	2.5×10^{-20}	$= \delta m_1^2$	5.6×10^{-18}	27.12	12	0.74%
Neutron dom. (1, 0, 0)	BPL	6.3×10^{-18}	1.0×10^{-21}	$= \delta m_2^2$	25.63	12	1.21%

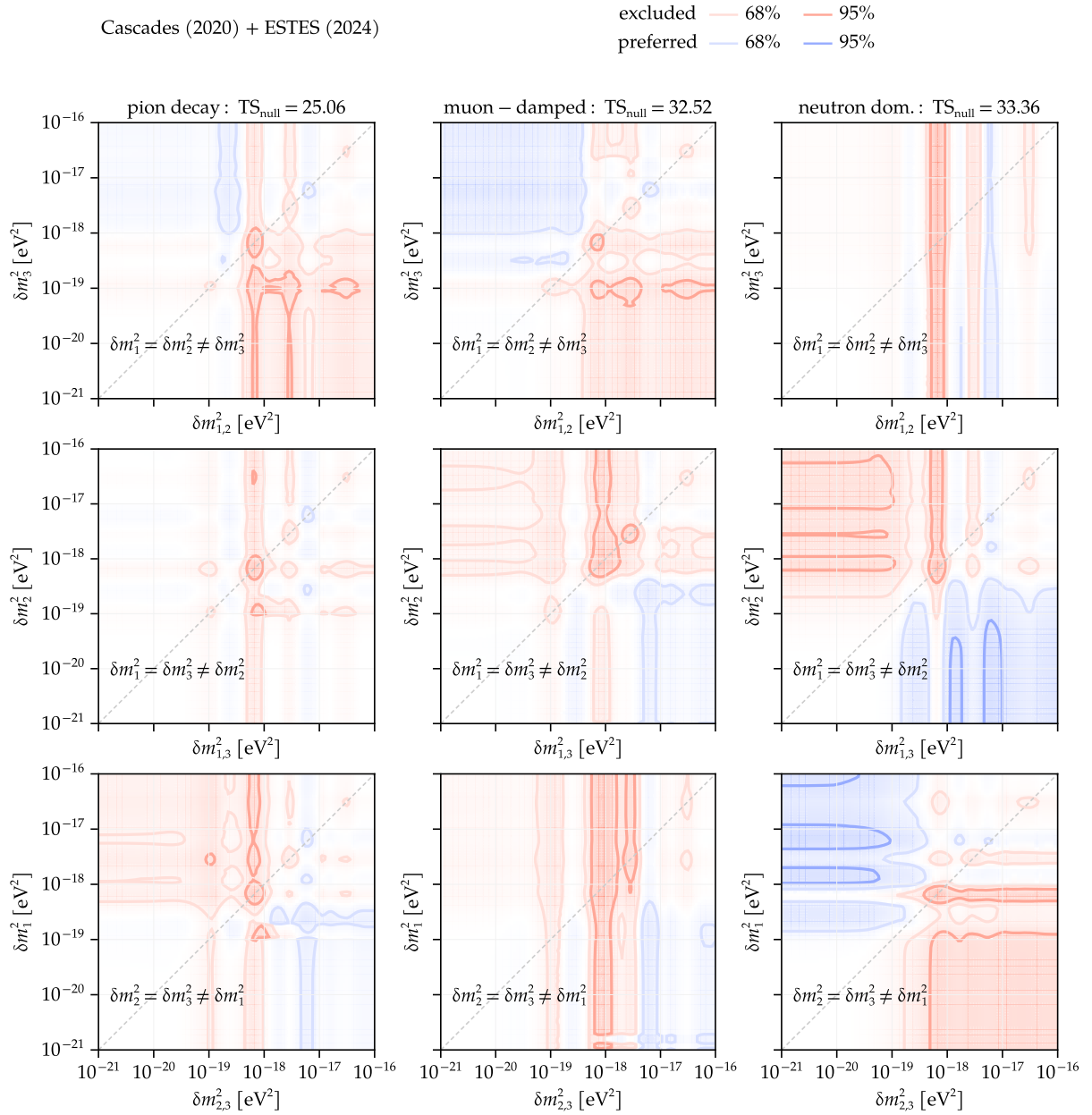


FIG. 7. Disfavored and preferred regions in two-dimensional quasi-Dirac parameter space. Left to right, the three columns correspond to the initial flavor ratios of the pion-decay, muon-damped, and neutron-dominated scenarios. These results use a source distribution over redshift that follows the SFRD, and assume the BPL emission spectrum.

Fig. 8: oscillations induced by $\delta m_1^2 = 2 \times 10^{-20} \text{ eV}^2$ disproportionately flatten the flux in the cascade channel at low energies, while oscillations induced by $\delta m_2^2 = \delta m_3^2 = 5.6 \times 10^{-18} \text{ eV}^2$ produce a dip in both the track and cascade channels at just below 1 PeV. Overall, the improvement to the fits in quasi-Dirac scenarios is not yet statistically significant, but provides an interesting alternative to flavor-uniform broken power-law models that can be further investigated in future works.

We also find that in the absence of quasi-Dirac effects, the muon-damped and neutron-dominated scenarios, which do not predict equal flavor ratios at Earth, produce significantly

worse fits than the pion-decay scenario. Quasi-Dirac models can slightly improve these fits: in the neutron-dominated case, a quasi-Dirac model with $\delta m_1^2 = 6.3 \times 10^{-18} \text{ eV}^2$, $\delta m_2^2 = \delta m_3^2 = 1 \times 10^{-21} \text{ eV}^2$ produces an overall test statistic comparable to the null pion-decay scenario, this model is illustrated in Fig. 9. A full grid of excluded and preferred regions for all three source flavor assumptions, assuming a BPL emission spectrum, is shown in Fig. 7.³

³For flavor-triangle analyses including quasi-Dirac neutrinos, see Refs. [32,33].

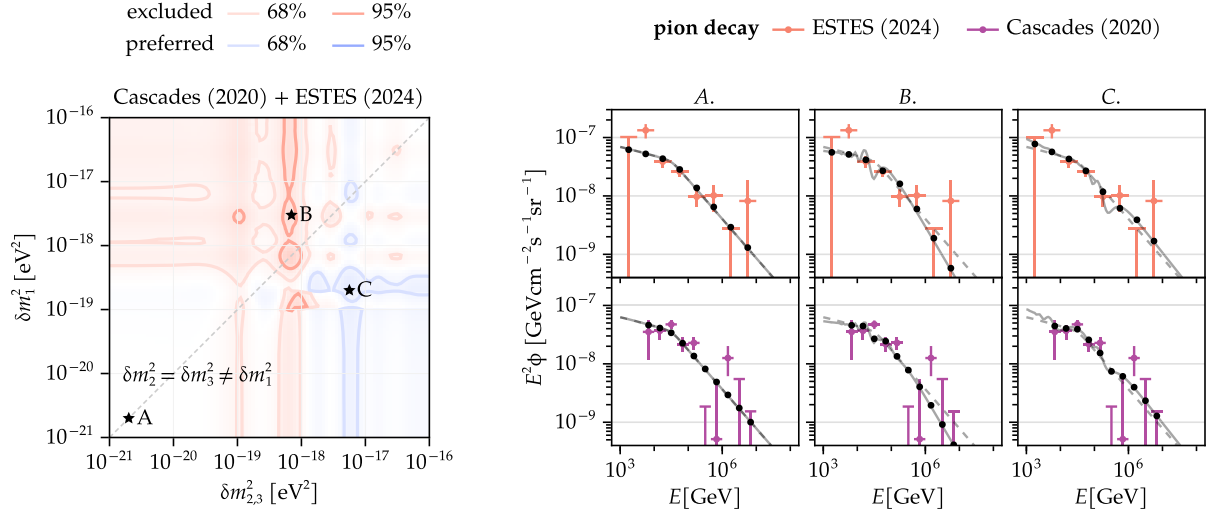


FIG. 8. Two-parameter quasi-Dirac models in the pion-decay production scenario. Left: the test statistic difference as a function of δm_1^2 and $\delta m_2^2 = \delta m_3^2$, assuming an initial flavor ratio produced by pion decays (1, 2, 0). Right: the best-fit quasi-Dirac flux hypotheses (solid gray curves) for the three points in two-dimensional quasi-Dirac parameter space are indicated, compared with the Cascades and ESTES piecewise fluxes. The black points indicate the predicted piecewise flux in each bin according to the quasi-Dirac model. The dashed gray curves indicate the null best fit.

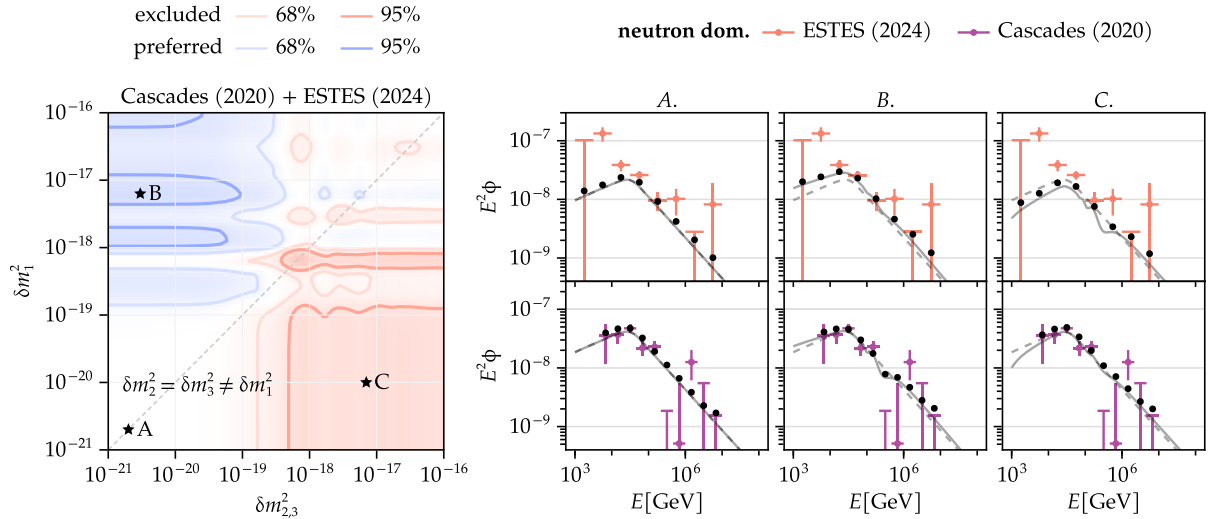


FIG. 9. Two-parameter quasi-Dirac models in the neutron-dominated production scenario. Left: the test statistic difference as a function of δm_1^2 and $\delta m_2^2 = \delta m_3^2$, assuming an initial flavor ratio dominated by neutron decays (1, 0, 0). Right: the best-fit quasi-Dirac flux hypotheses for the three points in two-dimensional quasi-Dirac parameter space are indicated, compared with the Cascades and ESTES piecewise fluxes.

VII. CONCLUSIONS

Neutrino telescopes provide a unique opportunity to uncover the origin of neutrino masses. In particular, the quasi-Dirac neutrino model, which is motivated by string theory scenarios, produces observable signatures in the spectra of astrophysical neutrinos. In this work, we have shown that quasi-Dirac oscillations remain observable from an extended population of sources. These produce observable disappearance signatures in the spectra and can impact the flavor ratio of astrophysical neutrinos.

Assuming that neutrinos are produced by a population of sources distributed over redshift z according to the SFRD, and that all sources emit neutrinos according to a broken power-law spectrum, we find that the most recent IceCube measurement of the astrophysical spectrum disfavors, at the 3σ level, a single quasi-Dirac mass-squared difference δm^2 between $(5 \times 10^{-19}, 8 \times 10^{-19})$ eV². Our sensitivity to this region is increased when we consider alternate possibilities for the source evolution in z , or alternate shapes for the emitted source spectrum.

Our excluded regions are robust across the combinations of spectral indices considered in this work, however, we also find an interesting interplay between the spectral description of the source flux and possible positive quasi-Dirac signatures. In particular, when the emitted spectrum is modeled either as a single unbroken power law or as a single power law with a spectral break, the data show a mild preference for quasi-Dirac neutrinos with mass splittings slightly below 10^{-19} eV². However, this scenario is not preferred over the standard no-quasi-Dirac solution with a broken power-law spectrum, reinforcing the overall conclusion of this work. Improved measurements below 100 TeV will be especially important for clarifying this situation and discriminating between these possibilities.

When we consider a model with two distinct mass splittings, which would produce energy-dependent modifications to the astrophysical flavor ratio, we find that quasi-Dirac models can explain the tension between track- and cascade-based measurements of the astrophysical spectrum, producing an overall better fit than a broken power-law model. These results should be further investigated when new track- and cascade-based measurements, which will leverage increased statistics and better systematic control, are published by IceCube and other neutrino telescopes.

The spectral and flavor distortion signatures produced by the quasi-Dirac scenario considered here may also arise in other new physics frameworks [69,70]. These include neutrino decay models [30,45,47,71–77], neutrino interactions with light dark matter [78–81], and effects induced by spacetime distortions or Lorentz-violating operators [82–86]. Among these possibilities, only neutrino decay can produce an L/E dependence similar to that of quasi-Dirac oscillations. Invisible decay scenarios, however, are already strongly constrained by cosmological observations [87–89] and flavor-based analyses [77]. Visible decays into different flavor combinations do not produce oscillatory signatures and are likewise constrained [77]. Light dark matter scenarios and Lorentz violation effects are also subject to existing flavor constraints [81,90,91]. If future data reveal evidence for quasi-Dirac-like features, dedicated analyses will be needed to assess how well these scenarios can be distinguished.

Future observations of the astrophysical flux, leveraging more years of data and reduced systematic uncertainties, will further improve our understanding of its

shape and flavor dependence, and thus further extend the sensitivity to neutrino mass models. Additionally, point source analyses are continuing to determine the origins of high-energy astrophysical neutrinos, constraining the properties of the source population and improving our sensitivity to propagation effects, such as quasi-Dirac-driven oscillations. These two beams are complementary. Hints of ultralong-baseline observations can be first observed by features in the diffuse, all-sky neutrino emission. Such signals could then be confirmed by the study of the neutrino yield from multiple sources.

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DATA AVAILABILITY

This study used publicly available results from IceCube Collaboration. This study's findings can be reproduced using the resources available at this Github repository: [92].

APPENDIX A: FLUX PARAMETER MARGINALIZATION

As mentioned in the main text, in our analyses we marginalize over the values of the parameters $\vec{\eta}$ of each source spectrum. For each parameter, we use a uniform prior, with allowed ranges as listed in Table IV.

TABLE IV. Allowed parameter ranges for flux parameters $\vec{\eta}$ of the SPL, SPE, and BPL flux models.

SPL		SPE		BPL	
$\log_{10} \phi_0$	$\in [-20, 20]$	$\log_{10} \phi_0$	$\in [-20, 20]$	$\log_{10} \phi_0$	$\in [-20, 20]$
γ	$\in [0, 4]$	γ	$\in [0, 4]$	γ_1	$\in [0, 4]$
				γ_2	$\in [1.5, 4]$
		$\log_{10} E_{\text{cut}}$	$\in [5, 10]$	$\log_{10} E_{\text{break}}$	$\in [4, 8]$

APPENDIX B: LIKELIHOOD CONSTRUCTION AND COMPATIBILITY WITH ICECUBE RESULTS

As noted in the main text, we approximate the IceCube likelihood function by the following function:

$$\text{TS}(\delta m_k^2) = \min_{\vec{\eta}} \sum_S \sum_i -2 \log \mathcal{L}_i(\mu_i^S(\delta m_k^2, \vec{\eta})). \quad (\text{B1})$$

This function needs to have the right statistical coverage, as reported by the IceCube results, while also having the proper limiting cases, which are distinct in signal- and background-dominated regimes. At sub-100 TeV energies, as the flux normalization goes to zero the likelihood should converge to a constant value, since this regime is background dominated. At supra-100 TeV energies, the data are dominated by astrophysical neutrinos, with negligible background contribution, and thus the flux normalization is proportional to the rate. This implies that in this regime the likelihood should behave asymptotically like a Poisson likelihood. Explicitly, for $E_\nu < 100$ TeV,

$$-2 \log \mathcal{L}_i(\mu_i^S(\delta m_k^2, \vec{\eta})) = \frac{(N_i^S - \mu_i^S(\delta m_k^2, \vec{\eta}))^2}{(\sigma_i^{\pm,S})^2}, \quad (\text{B2})$$

i.e., a Pearson χ^2 , and where $\sigma_i^{\pm,S}$ corresponds to the upper error ($\sigma_i^{+,S}$) on the data when $\mu_i^S > N_i^S$ and the lower errors ($\sigma_i^{-,S}$) otherwise. For $E_\nu > 100$ TeV,

$$\begin{aligned} & -2 \log \mathcal{L}_i(N_i^S | \mu_i^S(\delta m_k^2, \vec{\eta})) \\ &= -2 \left[\frac{\log \mathcal{L}^{\text{Poisson}}(N_i^S | \mu_i^S(\delta m_k^2, \vec{\eta})) - \log \mathcal{L}^{\text{Poisson}}(N_i^S | N_i^S)}{\log \mathcal{L}^{\text{Poisson}}(N_i^S | N_i^S \pm \sigma_i^{\pm,S}) - \log \mathcal{L}^{\text{Poisson}}(N_i^S | N_i^S)} \right], \end{aligned} \quad (\text{B3})$$

where $\mathcal{L}^{\text{Poisson}}(N|\mu)$ is the probability of observing N given expected μ , and the “+” in $\sigma_i^{\pm,S}$ is used for $\mu_i^S > N_i^S$ while the corresponding minus sign pair in the contrary case.

As can be seen in Table V, using this likelihood approximation, we reproduce all relevant results published by the experiment within uncertainties.

TABLE V. Compatibility with IceCube results. Best-fit parameters obtained in our analysis compared with IceCube’s published values and their uncertainties.

Fit	Sample	Flux model				
		SPL	γ	ϕ_0		
IceCube	Cascades		$2.53^{+0.07}_{-0.07}$	$1.66^{+0.25}_{-0.27}$		
This work	Cascades		2.41	1.95		
IceCube	ESTES		$2.58^{+0.1}_{-0.09}$	$1.68^{+0.19}_{-0.22}$		
This work	ESTES		2.57	1.60		
IceCube	CombinedFit		$2.52^{+0.036}_{-0.038}$	$1.8^{+0.13}_{-0.16}$		
This work	CombinedFit		2.35	1.83		
		SPE	γ	ϕ_0	$\log_{10} E_{\text{cut}}$	
IceCube	Cascades		$2.45^{+0.09}_{-0.11}$	$1.83^{+0.37}_{-0.31}$	$6.4^{+0.9}_{-0.4}$	
This work	Cascades		2.30	2.41	6.29	
IceCube	CombinedFit		$2.386^{+0.081}_{-0.09}$	$2.2^{+0.3}_{-0.25}$	$6.15^{+0.37}_{-0.24}$	
This work	CombinedFit		2.19	2.40	5.99	
		BPL	ϕ_0	$\log_{10} E_{\text{break}}$	γ_1	γ_2
IceCube	Cascades		$1.71^{+0.65}_{-0.29}$	$4.6^{+0.5}_{-0.2}$	$2.11^{+0.29}_{-0.67}$	$2.75^{+0.29}_{-0.14}$
This work	Cascades		2.15	4.42	1.52	2.67
IceCube	ESTES		$1.70^{+0.19}_{-0.22}$	4.36	$2.79^{+0.3}_{-0.5}$	$2.52^{+0.1}_{-0.09}$
This work	ESTES		1.51	4.67	2.46	2.72
IceCube	CombinedFit		$1.77^{+0.19}_{-0.18}$	$4.39^{+0.1}_{-0.1}$	$1.31^{+0.51}_{-1.3}$	$2.735^{+0.067}_{-0.075}$
This work	CombinedFit		1.75	4.36	1.25	2.74

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