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Third generation familons, B factories, and neutrino cosmology

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We study the physics of spontaneously broken family symmetries acting on the third generation. Massless familons (or Majorons) f associated with such broken symmetries are motivated especially by cosmological scenarios with decaying tau neutrinos. We first note that, in marked contrast to the case for the first two generations, constraints on third generation familon couplings are poor, and are, in fact, non-existent at present in the hadronic sector. We derive new bounds from $B^0-\overline{B^0}$ mixing, $B^0 \to l^+ l'^-$, $b \to s \nu \overline{\nu}$, and astrophysics. The resulting constraints on familon decay constants are still much weaker than those for the first and second generation. We then discuss the promising prospects for significant improvements from searches for $\tau \to l f$, $B \to (\pi, K) f$, and $b \to (d,s) f$ with the current CLEO, ARGUS, and CERN LEP data. Finally, we note that future constraints from CLEO III and the B factories will probe decay constants beyond 10^8 GeV, well within regions of parameter space favored by proposed scenarios in neutrino cosmology. [S0556-2821(98)03009-4]

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I. INTRODUCTION

For over half a century, one of the major puzzles in particle physics has been the question of why quark and lepton families replicate. Although we have accumulated a wealth of data concerning the masses and mixings of quarks and leptons, we still appear to be far from a true understanding of family structure. In the absence of a concrete model to consider, it is natural to postulate the existence of some family symmetry [1–3] that plays a role in determining the observed particle spectrum. Once we consider such a family symmetry, we face a plethora of options. The symmetry may be (1) discrete,[1] (2) continuous and local, or (3) continuous and global. Within each of these categories, one may choose any of a number of symmetry groups, and the overall family symmetry may even be a combination of the three possibilities.

Of course, any exact family symmetry of the underlying theory must be spontaneously broken at some energy scale since we know that the quark and lepton masses are very different from one family to the next. For option (1), spontaneously broken discrete symmetries, domain walls are the only model-independent predictions, and these cannot be studied in particle physics laboratories. In case (2), the masses of the family gauge bosons of spontaneously broken local continuous symmetries can be constrained, e.g., from K^0-$\overline{K}^0$ mixing [5].

From a phenomenological point of view, however, possibility (3) is particularly enticing, as it implies the existence of massless Nambu-Goldstone bosons, called ‘‘familons,’’ from the spontaneously broken family symmetry. This family symmetry may be either Abelian or non-Abelian; Nambu-Goldstone bosons associated with the spontaneous symmetry breaking of an Abelian lepton number symmetry are often called ‘‘Majorons.’’ [2] The existence of new massless particles has many implications in particle physics, astrophysics, and cosmology, and, as we will see, may be probed in a wide variety of experiments. Moreover, the couplings of familons at low energies are determined by the non-linear realization of the family symmetry. These couplings are, e.g., of the form

$$\frac{1}{F} \partial_\mu f^a \overline{\psi_L^i} \gamma^\mu T^a_{ij} \psi_L^j , \tag{1.1}$$

where F is the family symmetry breaking scale, i.e., the familon decay constant, f^a are the familons, T^a are the generators of the broken symmetry, and the ψ_L are fermion fields in terms of which the flavor symmetry is defined. The strength of the familon coupling is therefore inversely proportional to F and can be constrained for a given family symmetry group in a model-independent manner.

Familon couplings between the first and second generations have been studied extensively and will be reviewed below. In contrast, however, couplings involving the third generation are largely unexplored, although they may have rather rich phenomenological and cosmological implications [8]. Current constraints in the lepton sector are relatively weak, with the best bounds coming from $\tau \to (e, \mu) f$ bounds [9], and there are at present no corresponding bounds reported in the hadronic sector (see, however, Ref. [10]). At the same time, it is a logical possibility that the familon

[1]There is a subtle distinction between global and gauged discrete symmetries [4]. For this phenomenological analysis, however, they are equivalent.

[2]Majorons have been extensively studied, and arise in a variety of models [6], including, for example, supersymmetric theories with spontaneous R-parity breaking [7]. In this paper, we study a number of probes, many of which are applicable to both Abelian and non-Abelian symmetries. We use the generic name ‘‘familon’’ to denote the associated Nambu-Goldstone bosons in either case.

couples preferentially to the third generation, and models have been proposed in which this is the case [11]. It is therefore interesting to explore the possibilities for improving (or setting) bounds on familon scales for the third generation, especially in light of the upcoming B physics experiments.

In this paper, we will study what we believe to be the most sensitive probes of couplings of familons to the third generation, primarily to τ leptons and b quarks. We show that dedicated analyses of existent data from CLEO, ARGUS, and the CERN e^+e^- collider LEP could probe family symmetry breaking scales up to $\sim 10^7$ GeV and may be significantly improved at future B factories. Simply because this is largely unexplored physics, there is a high discovery potential for familons at these facilities.

Familon couplings to the third generation are also of interest from a cosmological point of view. The mass of the τ neutrino is still allowed to be as large as 18.2 MeV experimentally [12]. A heavy τ neutrino has interesting consequences for both big-bang nucleosynthesis (BBN) [13–16] and large scale structure formation [17–20], as will be discussed in Sec. VI. Since a heavy neutrino ($\geqslant 100h^2$ eV, where h is the expansion rate of the universe in units of 100 km/sec/Mpc) must decay in order not to overclose the universe, an invisible decay into a lighter neutrino and a massless boson, such as a familon (or Majoron), is typically required. (The three neutrino mode is strongly disfavored and therefore the familon mode is most preferred [13].) There is therefore an interesting interplay between experimental searches for familons and scenarios requiring heavy neutrinos, and, as we will see, future collider experiments and analyses may severely constrain a number of such cosmologically motivated scenarios.

This paper is organized as follows. We begin in Sec. II with a discussion of familon interactions. In particular, we emphasize that the familon interactions of particles in the same gauge multiplet are expected to be comparable. In Sec. III we consider constraints on familon interactions that may be inferred from current experimental data, concentrating on familon couplings to the third generation. Current bounds on third generation couplings from astrophysical considerations are presented in Sec. IV. We then describe some promising prospects for detecting familons in B physics at future experiments in Sec. V. Finally, we note some of the interesting cosmological implications in Sec. VI and give our conclusions in Sec. VII.

II. FAMILON INTERACTIONS

The standard model contains 15 particle states in each of the 3 generations. These states are distinguished by the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge interactions, which divide each generation into 5 multiplets: Q, U, D, L, and E. The gauge interactions therefore break the flavor symmetry group from U(45) to $U(3)^5$. In the standard model, the flavor group $U(3)^5$ is broken explicitly to $U(1)_B \times U(1)_L$ by Yukawa couplings. However, in extensions of the standard model in which one hopes to gain some understanding of the pattern of fermion masses and mixings, some subgroup of the flavor group may be an exact symmetry of the Lagrangian that is broken spontaneously by the vacuum, and it is this possibility we consider here.

The massless Nambu-Goldstone bosons of the spontaneously broken flavor symmetry, familons [1–3], have interactions given by the couplings

$$\mathcal{L}_f = \frac{1}{F}\partial_\mu f^a J^{\mu a}, \tag{2.1}$$

where f^a are the familon fields, and $J^{\mu a}$ are flavor currents.[3] The interactions are suppressed by F, the scale at which the flavor symmetry is spontaneously broken. Note that familons are derivatively coupled,[4] and so do not mediate long-range ($\sim r^{-2}$) forces. The most general current $J^{\mu a}$ composed of two fermion fields takes either the form

$$J^{\mu a} = \overline{\psi_i}\gamma^\mu (g_V + g_A\gamma_5) T^a_{ij}\psi_j \tag{2.2}$$

or

$$J^{\mu a} = \overline{\psi_i}\gamma^\mu (g_L P_L + g_R P_R) T^a_{ij}\psi_j , \tag{2.3}$$

where $P_{L,R}$ are the projection operators $\frac{1}{2}(1\pm\gamma_5)$, i and j are generational indices, and T^a_{ij} are the spontaneously broken generators of the family symmetry. The fields ψ_i and ψ_j are fermion mass eigenstates, which we assume here to be also flavor eigenstates. (The more general case is described below.) Using the form of the current given in Eq. (2.2), the familon interaction may be written as

$$\begin{aligned}\mathcal{L}_f &= \frac{1}{F}\partial_\mu f^a \overline{\psi_i}\gamma^\mu (g_V + g_A\gamma_5) T^a_{ij}\psi_j \\ &= -\frac{i}{F} f^a \overline{\psi_i}[g_V(m_i - m_j) + g_A(m_i + m_j)\gamma_5] T^a_{ij}\psi_j ,\end{aligned} \tag{2.4}$$

where in the last step we have integrated by parts and then substituted the equations of motion. The second line of Eq. (2.4) is of course only valid for on-shell fermions such as external leptons, whereas in hadronic matrix elements and processes including off-shell fermions, the derivative coupling of the first line must be used.

We see that familons may mediate or be produced in family-changing processes. They may also couple to identical fermions $\psi_i = \psi_j$, but only through axial couplings. What processes are mediated by familons depends on the particular family symmetry group that is broken. For example, for O(N) groups, the generators T_{ij} are anti-symmetric, and so do not generate flavor-diagonal interactions. However, they do generate interactions like $f\overline{\psi_i}\gamma_5\psi_j - f\overline{\psi_j}\gamma_5\psi_i$, where we have considered axial vector current interactions as an example. Familons from O(N) groups may therefore mediate neutral meson mixing, which we will consider in Sec. III B. The situation is reversed for SU(N) groups. Here, flavor-diagonal couplings exist. However, if we consider any SU(2)

[3]Throughout this study, we will assume that no additional light degrees of freedom are introduced by other new physics.

[4]If the flavor symmetry is anomalous, familons may also have non-derivative, flavor-diagonal couplings. We will not consider such couplings here.

subgroup and form the complex familon $\tilde{f}=f^1+if^2$, the off-diagonal interactions are given by $\tilde{f}\overline{\psi_i}\gamma_5\psi_j+\tilde{f}^*\overline{\psi_j}\gamma_5\psi_i$, and we see that $\tilde{f}$ exchange cannot induce neutral meson mixing.

Up to this point, we have ignored possible mass mixing effects. In general, if the flavor eigenstates ψ' are related to the mass eigenstates ψ by

$$\psi'=U_\psi\psi, \tag{2.5}$$

where U_ψ is a 3×3 unitary mixing matrix, the familon interactions are given by

$$\mathcal{L}_f=\frac{1}{F}\partial_\mu f^a\overline{\psi_i'}\gamma^\mu(g_V+g_A\gamma_5)T^a_{ij}\psi_j'$$
$$=\frac{1}{F}\partial_\mu f^a\overline{\psi_i}\gamma^\mu(g_V+g_A\gamma_5)T^a_{\psi ij}\psi_j, \tag{2.6}$$

where $T^a_\psi=U^\dagger_\psi T^a U_\psi$. Mass mixings may therefore generate flavor-diagonal interactions from flavor off-diagonal interactions, and vice versa. For example, in the case of an Abelian U(1) symmetry, mass mixing effects may generate flavor-changing interactions. They may also extend non-maximal family symmetries to couplings involving all three generations; for example, a U(2) symmetry between the first and second families, may, after rotation to mass eigenstates, result in familon interactions involving the third generation.

While the phenomenology of familons varies from group to group, it is important to note that gauge symmetry relates the familon interactions of particles in the same gauge multiplet. As an example, let us consider a spontaneously broken lepton flavor symmetry. The familon interaction is then given by

$$\frac{g_L}{F}\partial_\mu f^a\overline{L_i'}\gamma^\mu T^a_{ij}L_j', \tag{2.7}$$

where the SU(2) lepton doublets $L_i'=(\nu_i',l_i')$ are in the flavor eigenstate basis. This interaction therefore generates familon interactions for both the charged leptons and neutrinos. In the presence of neutrino masses, the flavor eigenstates may not correspond to mass eigenstates. The familon interactions in the mass basis are then

$$\frac{g_L}{F}\partial_\mu f^a\overline{\nu_i}\gamma^\mu T^a_{\nu ij}\nu_j+\frac{g_L}{F}\partial_\mu f^a\overline{l}_i\gamma^\mu T^a_{lij}l_j, \tag{2.8}$$

where $T^a_\nu=U^\dagger_\nu T^a U_\nu$, and we have defined $V=U^\dagger_\nu U_l$ and $T^a_l=V^\dagger T^a_\nu V$. T^a_l and T^a_ν are therefore related by a similarity transformation, and in the presence of mass mixing, the couplings of the interactions of $\partial_\mu f^a\overline{\nu_i}\gamma^\mu\nu_j$ and $\partial_\mu f^a\overline{l}_i\gamma^\mu l_j$ are not necessarily identical. However, in the absence of fine-tuning, we expect these couplings to be of the same magnitude. Bounds on one familon interaction may thus be considered to imply comparable bounds on the other interactions linked by gauge symmetry.

Because the familon interactions of particles in the same gauge multiplet are comparable in the absence of fine-tuning, there are many more relations in theories with enlarged gauge groups. For example, for SU(5) grand unified theories (GUTs), the particles $d_R,\nu,e_L\subset\overline{\mathbf{5}}$ are expected to have comparable familon interactions, as are the particles $u_L,d_L,u_R,e_R\subset\mathbf{10}$. A particularly relevant example for our study below is that, in the GUT framework, bounds on familon decays of B mesons imply bounds on familon decays of tau neutrinos in the absence of fine-tuning.

Flavor mixing effects also induce familon couplings of fields with different generational indices. In the quark sector, for example, substituting the quark doublet $Q_i'=(u_i',d_i')$ for L_i' in the discussion above, Eq. (2.8) becomes

$$\frac{g_L}{F}\partial_\mu f^a\overline{u}_i\gamma^\mu T^a_{uij}u_j+\frac{g_L}{F}\partial_\mu f^a\overline{d}_i\gamma^\mu T^a_{dij}d_j, \tag{2.9}$$

where T^a_u and T^a_d are related by the Cabibbo-Kobayashi-Maskawa (CKM) matrix through

$$T^a_d=V^\dagger_{\text{CKM}}T^a_u V_{\text{CKM}}. \tag{2.10}$$

We see that in general, couplings to all generations are induced by flavor mixings. For example, a familon with flavor-diagonal coupling to $\overline{t}t$ in the up sector couples not only to $\overline{b}b$, but also to, for example, $\overline{b}s$ and $\overline{d}d$. The induced couplings to first and second generation quarks in this case are CKM-suppressed, but may still lead to significant bounds when, as is often the case, these induced couplings are much more strongly constrained. We will consider the constraints on mixing-induced couplings from K decays in Sec. III A and from supernova cooling in Sec. IV C.

Finally, note in Eq. (2.2) that the strength of the interaction depends not only on F, but also on T^a_{ij} and the couplings $g_{V,A}$. In the following sections, we will present a variety of bounds on combinations of these couplings, and it is important that we define our conventions and normalizations. We will always define our interaction as

$$\frac{1}{F}\partial_\mu f\overline{\psi_i}\gamma^\mu(g^{ij}_V+g^{ij}_A\gamma_5)\psi_j, \tag{2.11}$$

and similarly for g^{ij}_L and g^{ij}_R; the superscripts of the couplings will often be omitted when they are obvious from the context. In presenting our bounds, it will be convenient to define

$$F^I_{ij}\equiv F/g^{ij}_I, \tag{2.12}$$

where $I=V,A,L,R$. In addition, as many of our bounds are to a good approximation independent of the chirality of the interaction and so only dependent on the combination $g^{ij2}_V+g^{ij2}_A$, we define

$$F_{ij}\equiv\frac{F}{\sqrt{g^{ij2}_V+g^{ij2}_A}}. \tag{2.13}$$

III. BOUNDS FROM ACCELERATOR DATA

As described in the previous section, familons may take part in flavor-changing processes, and bounds on such processes lead to lower bounds on the familon energy scale. For familons mediating transitions between the first and second generation, such bounds are rather stringent. In contrast, similar bounds involving the third generation are much weaker, with the previously reported constraints limited only to bounds from rare τ decays. We are thus motivated to focus on the third generation. In Sec. III A, we begin by reviewing and contrasting such bounds, and then discuss the implications of flavor eigenstate mixings. We then go on to derive new bounds from a variety of processes. In Sec. III B we consider familon-mediated processes such as neutral meson mixing and rare leptonic decays of mesons. Finally, in Sec. III C we consider possible analyses at LEP and extrapolate a preliminary ALEPH bound on $b\to s\nu\bar{\nu}$ to a bound on $b\to sf$.

A. Decays to familons

We begin by considering bounds from decays of mesons and leptons to familons. Normalizing the relevant familon scale according to Eq. (2.11), we find

$$\Gamma(K\to\pi f)=\frac{1}{16\pi}\frac{m_K^3}{F^2}g_V^2\beta^3|F_1(0)|^2, \tag{3.1}$$

where $\beta=1-m_\pi^2/m_K^2$. In the limit of exact flavor SU(3) symmetry, the form factor $\langle\pi^+(p')|\bar{s}\gamma^\mu d|K^+(p)\rangle=F_1(q^2)(p+p')^\mu$ at zero momentum transfer has a fixed normalization, $F_1(0)=1$. For leptonic decays $l_i\to l_j f$, the exact tree-level partial decay width in the limit of massless l_j is given by

$$\Gamma(l_i^-\to l_j^- f)=\frac{1}{16\pi}\frac{m_{l_i}^3}{F^2}(g_V^2+g_A^2)\beta^3, \tag{3.2}$$

where here $\beta=1-m_{l_j}^2/m_{l_i}^2$.

The strongest bound on any flavor scale is derived from the constraint on exotic K decay. Using the above expressions, the experimental result $B(K^+\to\pi^+ f)<3.0\times10^{-10}$(90% C.L.) [21] leads to the bound

$$F_{sd}^V>3.4\times10^{11}\ \text{GeV}. \tag{3.3}$$

Note that the limit on $B(K^+\to\pi^+ f)$ bounds only the vectorial familon coupling; the axial coupling is unconstrained. For the leptonic sector, Jodidio *et al.* report the constraint $B(\mu^+\to e^+ f)<2.6\times10^{-6}$(90% C.L.) [22], which they obtain under the assumption of a vector-like familon coupling. This can be converted into the bound

$$F_{\mu e}^V>5.5\times10^9\ \text{GeV}. \tag{3.4}$$

For familon interactions of arbitrary chirality, the slightly weaker constraint

$$F_{\mu e}>3.1\times10^9\ \text{GeV} \tag{3.5}$$

may be obtained from the bound $B(\mu^+\to e^+\gamma f)<1.1\times10^{-9}$(90% C.L.) [23].

We now compare these bounds to those available in the third generation. The ARGUS Collaboration [9] has bounded the branching fractions of τ decays into light bosons and found the limits $B(\tau^-\to\mu^- f)<4.6\times10^{-3}$(95% C.L.) and $B(\tau^-\to e^- f)<2.6\times10^{-3}$(95% C.L.). These imply the following constraints on the flavor scale:

$$F_{\tau\mu}>3.2\times10^6\ \text{GeV} \tag{3.6}$$

$$F_{\tau e}>4.4\times10^6\ \text{GeV}. \tag{3.7}$$

We see that the bounds on flavor scales in the leptonic sector are significantly less stringent for third generation couplings than for those involving only the first two. The discrepancy is even more pronounced in the hadronic sector, where there are as yet no reported bounds on flavor scales from B decays.

It is also worth noting, however, that strong bounds on a particular flavor scale, such as the one on F_{sd}^V, may imply significant bounds on other flavor scales as well. These bounds are induced by the flavor-mixing effects discussed in Sec. II and are thus model-dependent. As an example let us now assume that flavor and mass eigenstates coincide for up-type quarks. A given familon coupling in the up sector requires, by gauge invariance, a corresponding coupling in the down sector. For example, from Eqs. (2.9) and (2.10) we see that the coupling $\partial_\mu f\bar{t}\gamma^\mu P_L c/F_{tc}^L$ induces the coupling $V_{ts}^*V_{cd}\partial_\mu f\bar{s}\gamma^\mu P_L d/F_{tc}^L$, which mediates the rare decay $K^+\to\pi^+ f$. Assuming complex familons, the Hermitian conjugate coupling gives a similar contribution $\propto V_{cs}^*V_{td}$ to the decay into the complex conjugate familon. Summing both decay widths and comparing to the bound on F_{sd}^V in Eq. (3.3), one can derive the mixing induced bound

$$F_{tc}^L>2.2\times10^9\ \text{GeV}. \tag{3.8}$$

Under similar assumptions, we find $F_{tu}^L>6.6\times10^9$ GeV. Note, however, that such bounds do not apply if the mass and flavor bases are aligned in the down sector [24] or if the couplings are purely axial.

B. Familon-mediated processes

In this section we derive new constraints on the scale of spontaneous flavor symmetry breaking by considering nonstandard familon contributions to neutral meson mixing and existing bounds on rare leptonic decays such as $B^0\to\tau e$.

A familon contribution to neutral meson mixing requires a real flavor group to be spontaneously broken in the corresponding sector, such that the same real familon scalar field couples to the quark current and its Hermitian conjugate. For concreteness, let us consider the B^0-$\overline{B^0}$ system; similar formulae hold (at least approximately) for other neutral meson systems. Assuming the general coupling structure

$$\frac{i}{F}\partial_\mu f[\bar{d}\gamma^\mu(g_V+g_A\gamma_5)b-\bar{b}\gamma^\mu(g_V+g_A\gamma_5)d], \tag{3.9}$$

we find a familon contribution to the mass splitting of

$$\Delta m_{B^0}^{(f)}\equiv|m_{B^0}-m_{\overline{B^0}}|\approx\frac{5}{6}\frac{f_{B^0}^2 g_A^2 m_{B^0}}{F^2}. \tag{3.10}$$

Equation (3.10) may be derived by taking the matrix element of the non-local operator

$$\frac{1}{2!}\frac{1}{F^2}\overline{d}_\alpha\gamma^\mu(g_V+g_A\gamma_5)b_\beta\frac{q_\mu q_\nu}{q^2}\overline{d}_\gamma\gamma^\nu(g_V+g_A\gamma_5)b_\delta \tag{3.11}$$

between B^0 and $\overline{B^0}$ states and using the definition of the pseudoscalar decay constant, $\langle 0|\overline{b}\gamma^\mu\gamma_5 d(0)|B^0(p)\rangle = if_{B^0}p^\mu$. The subscripts α, β, γ, and δ in Eq. (3.11) are color indices. Between two color singlet states, there are two contributions. The first one arises from $\alpha=\beta$ and $\gamma=\delta$ with a familon in the s-channel. In this case, the momentum transfer through the familon propagator is $q^2=m_{B^0}^2$, and after a vacuum insertion, it is easy to verify that this contribution is as in Eq. (3.10), but without the factor of 5/6. However, there is also a t-channel contribution from $\alpha=\delta$ and $\beta=\gamma$, which may be evaluated by a Fierz transformation and then a vacuum insertion as before. For a heavy–light system like the B^0 meson, one may assume the free-quark picture, in which the momentum transfer is governed by the energy of the ''static'' b quark $q^0\approx m_b\approx m_{B^0}$, and, in the numerator, the derivative acting on the quark current gives again a factor of m_b. Using $\langle 0|\overline{b}\gamma_5 d(0)|B^0(p)\rangle\approx if_{B^0}m_{B^0}$ and including the relative color factor of 1/3, one can estimate the t-channel contribution to be $-1/6$ times the s-channel contribution, which leads to Eq. (3.10).

Our result should be fairly reliable for the B^0 meson. For D^0 and K^0 mesons, the evaluation of the t-channel momentum transfer is more ambiguous. However, because this contribution is suppressed relative to the s-channel part, we expect the result of Eq. (3.10) to be reasonably accurate in these cases as well. We also note that a vector-like familon interaction does not contribute to the mass splitting, at least in the heavy quark approximation $m_b\approx m_{B^0}$. Although one might expect a vector contribution to appear in the t-channel contribution after the Fierz rearrangement, one finds that the term proportional to g_V contains axial vector and pseudoscalar contributions of equal magnitude but opposite sign.

The constraint on the flavor scale F results in principle from the requirement that the combined standard model and familon contributions do not exceed the measured value. However, when considering nonstandard contributions, it is also uncertain what one should take as the standard model contribution. For example, the reported value [25] for $|V_{tb}^*V_{td}|$ is derived from B^0-$\overline{B^0}$ mixing under the assumption that the standard model gives the only contribution. As a conservative bound, we simply compare the familon contributions directly to the corresponding measured values. The results are summarized in Table I. We take the decay constants to be $f_{B^0}\approx 175$ MeV and $f_{D^0}\approx 205$ MeV from recent lattice results [26], and $f_{K^0}\approx f_{K^+}\approx 160$ MeV [25]. Since we use the measured mass splitting (not its error), the bounds from B^0 and K^0 will only improve when the size of the standard model contribution can be quantified independently. For the D^0, where only the upper bound on the mass splitting is known, future experiments will improve the bound.

We next consider rare leptonic decays of neutral mesons, mediated by familon exchange. Such decays are possible if the same familon couples to both quarks and leptons. This is guaranteed in grand unified scenarios, where quarks and leptons are in the same gauge multiplet. In general the relevant interaction can be written in terms of effective vector and axial vector couplings that parametrize the familon couplings and mixing angles of a particular model. For example, the process $B^0\to\tau^+e^-$ can be mediated by the interaction Lagrangian

$$\frac{1}{F}\partial_\mu f[\overline{b}\gamma^\mu(g_V^{bd}+g_A^{bd}\gamma_5)d+\overline{\tau}\gamma^\mu(g_V^{\tau e}+g_A^{\tau e}\gamma_5)e]+\text{H.c.} \tag{3.12}$$

Note that the constants g_V and g_A may be different in the hadronic and leptonic sectors. Also, even if familon couplings always include third generation flavor eigenstates, mixing effects may induce transitions such as $B^0\to\mu^+e^-$.

With the interaction defined in Eq. (3.12) one obtains a width of

$$\Gamma(B^0\to\tau^+e^-)\approx\frac{1}{8\pi}\frac{f_{B^0}^2 g_A^{bd2}m_{B^0}m_\tau^2}{F^4}\left[(g_V^{\tau e2}+g_A^{\tau e2})\beta^2 - 2\frac{m_e}{m_\tau}(g_V^{\tau e2}-g_A^{\tau e2})\beta\right], \tag{3.13}$$

where $\beta=1-m_\tau^2/m_{B^0}^2$, and we have displayed the leading $g_V^2-g_A^2$ piece. In the limit where the lighter lepton is massless, the result is independent of the chirality of the interaction and depends only on the combination of lepton couplings $g_V^2+g_A^2$. Expressions for other similar processes are obtained by replacing the coupling constants g_V, g_A and the meson and lepton masses accordingly. Limits on the flavor scales from current experimental bounds on rare leptonic decays are given in Table II.

TABLE I. Bounds on the flavor scale from contributions to neutral meson mixing from familon exchange as given in Eq. (3.10). Note that these limits do not apply to vector-like couplings, and that this process requires a real flavor group so that a real familon scalar field couples to a current operator and its Hermitian conjugate, as in Eq. (3.9).

	$\Delta m_{\rm exp}$	Bound
$B^0-\overline{B^0}$	$0.5\times 10^{12}\ \hbar s^{-1}$ [25]	$F_{bd}^A>6.4\times 10^5$ GeV
$D^0-\overline{D^0}$	$<21\times 10^{10}\ \hbar s^{-1}$ [27]	$F_{cu}^A>6.9\times 10^5$ GeV
$K^0-\overline{K^0}$	$0.53\times 10^{10}\ \hbar s^{-1}$ [25]	$F_{sd}^A>1.7\times 10^6$ GeV

TABLE II. Limits on flavor scales and couplings for some rare meson decays. The branching ratio bounds are on the sum of the two charge states, assuming real familon scalars that mediate both decay modes. If familons mediate only one decay mode, the quoted bounds on F are weakened by a factor of $2^{1/4}$. In calculating these bounds, we neglect small corrections from the lighter lepton mass.

	Branching ratio upper bound	Bound
$B^0\to\tau^\pm e^\mp$	5.3×10^{-4}[28]	$(F_{bd}^A F_{\tau e})^{1/2}>3.5\times 10^3$ GeV
$B^0\to\tau^\pm\mu^\mp$	8.3×10^{-4}[28]	$(F_{bd}^A F_{\tau\mu})^{1/2}>3.1\times 10^3$ GeV
$B^0\to\mu^\pm e^\mp$	5.9×10^{-6}[28]	$(F_{bd}^A F_{\mu e})^{1/2}>2.8\times 10^3$ GeV
$D^0\to\mu^\pm e^\mp$	1.9×10^{-5}[29]	$(F_{cu}^A F_{\mu e})^{1/2}>1.2\times 10^3$ GeV
$K_L^0\to\mu^\pm e^\mp$	3.3×10^{-11}[30]	$(F_{sd}^A F_{\mu e})^{1/2}>3.8\times 10^5$ GeV

The bounds of Tables I and II are significantly weaker than those presented in Sec. III A. This is especially true in Table II, as rare leptonic meson decays are dependent on the flavor scale to the fourth power. However, such processes set bounds on third generation hadronic familon couplings, which were previously unconstrained. It is also important to note that the bounds on familon couplings to the first two generations are also interesting, as they constrain axial couplings, whereas the bound from K decay reviewed in the previous section bounds only vector-like couplings.

C. Constraints from LEP

Currently, there are no reported experimental bounds on decays $b \to (s,d)f$. One can, however, infer a constraint from ALEPH's preliminary bound on $b \to s\nu\bar{\nu}$ [31]. By searching for events with large missing energy, they placed the constraint $B(b \to s\nu\bar{\nu}) < 7.7 \times 10^{-4}$(90% C.L.). One can rescale this constraint to obtain an upper bound on $B(b \to sf)$.

The analysis for $b \to s\nu\bar{\nu}$ relies on the $E_{\rm miss}$ distribution [32], where $E_{\rm miss}$ is defined by

$$E_{\rm miss} = E_{\rm beam} + E_{\rm corr} - E_{\rm vis} \tag{3.14}$$

in each hemisphere of b-tagged events. Here, $E_{\rm beam}$ is half of the center-of-momentum energy, $E_{\rm corr} = (M^2_{\rm same} - M^2_{\rm opp})/4E_{\rm beam}$, where $M_{\rm same}$ and $M_{\rm opp}$ are the visible invariant masses in the same and opposite hemispheres, respectively, and $E_{\rm vis}$ is the total visible energy in the hemisphere. $E_{\rm corr}$ improves the estimate of the actual missing energy in the hemisphere by correcting for the fact that the hemisphere with larger invariant mass typically has higher energy.

The backgrounds from $b \to l\nu X$ and $c \to l\nu X$ are suppressed by rejecting events with identified $e^{\pm}$ or $\mu^{\pm}$ in the relevant hemisphere. Up to this point, we do not expect significant differences in efficiencies between the $b \to s\nu\bar{\nu}$ mode and the $b \to sf$ mode. They then required 35 GeV$< E_{\rm miss} <$45 GeV. The efficiencies for this requirement obviously differ between the two decay modes, since the mode $b \to s\nu\bar{\nu}$ has two missing neutrinos, resulting in a harder $E_{\rm miss}$ spectrum than that of the $b \to sf$ mode. The $E_{\rm miss}$ spectrum of both modes may be calculated by convoluting the theoretical missing energy distribution in three-body ($s\nu\bar{\nu}$) and two-body (sf) decays with the measured b fragmentation function [33]. We find that the ratio of efficiencies is 0.43 with little dependence on the details of the fragmentation function. By scaling the reported $B(b \to s\nu\bar{\nu})$ upper bound by this factor, we find

$$B(b \to sf) < 1.8 \times 10^{-3}. \tag{3.15}$$

Using the expression of Eq. (3.2) with the substitution of $m_{B^0} \approx m_b$ for m_{l_i}, this corresponds to a limit on the flavor scale of

$$F_{bs} > 6.1 \times 10^7 \ \text{GeV}. \tag{3.16}$$

Note that this analysis does not require an energetic strange particle, and so the constraint of Eq. (3.15) is actually on the sum $B(b \to sf) + B(b \to df)$. Thus, for $F_{bs} \approx F_{bd}$, the bound on the flavor scale given in Eq. (3.16) improves by a factor $\sqrt{2}$. The bound of Eq. (3.16) is enhanced by the fact that the SM decay width is greatly suppressed by V_{cb}, which increases the sensitivity of b decays to small exotic decay widths.

IV. BOUNDS FROM ASTROPHYSICS

In this section, we discuss constraints on third generation familon couplings from astrophysics. We begin in Sec. IV A with constraints on direct (tree-level) couplings. Second and third generation particles are absent in almost all astrophysical objects. The exception is supernovae, where all three neutrino species are thermalized in the core. We therefore consider what bounds on familon couplings to τ neutrinos may be obtained by supernova observations. Couplings of familons to the third generation may also radiatively induce couplings to first generation particles. Although such induced couplings are suppressed by loop factors, they are so stringently bounded by constraints from supernovae, white dwarfs, and red giants that interesting bounds also result. These are studied in Sec. IV B. Finally, mixings of flavor eigenstates may also induce couplings of familons to the first generation; such effects are discussed in Sec. IV C. It is important to note that, while the bounds derived in this section are rather strong in certain cases, they are also typically more model-dependent than, for example, the accelerator bounds of the previous section. We therefore specify the necessary conditions for each bound in detail in each case.

A. Bounds from direct couplings

In 1987, the Kamiokande group and the IMB group independently detected neutrinos emitted from supernova SN 1987A. They observed that the neutrino pulse lasted for a few seconds. Furthermore, their results indicate that neutrinos carried off about 10^{53} erg from the supernova. The observed duration time and neutrino flux can be well explained by the generally accepted theory of core collapse, and the observations confirmed the idea that most of the released energy in the cooling process is carried off by neutrinos. Exotic light particles, such as familons, may affect the agreement of theory and observation, since they can also carry off a significant energy fraction. The core of the supernova is hot ($T \sim 30$ MeV) and dense, and so neutrinos are thermalized in the core and can be a source of familon emission. If the energy fraction carried away by familons is substantial, the duration time of the neutrino pulse becomes much shorter than the observed value. In order not to affect the standard cooling process, the familon luminosity Q_f must be smaller than the neutrino luminosity, i.e., less than $\sim 10^{53}$ erg/sec.

This constraint can be satisfied in two different regimes of the familon coupling strength. For sufficiently high flavor scales F, the familon interaction is weak enough that familons are rarely produced and the familon luminosity Q_f is suppressed. On the other hand, for sufficiently low flavor scales, although familons are readily produced, they interact so strongly that they become thermalized and trapped in the core as well, thus decreasing the familon luminosity. Therefore, there are two parameter regions consistent with observations, high and low F, with an excluded region in the middle.

These types of constraints have been discussed by Choi and Santamaria [34] in the context of a Majoron model. We modify their discussions slightly for the familon case. To simplify the analysis, we will look at two extreme scenarios. First we consider a diagonal familon coupling to ν_τ, as in the case of an Abelian family symmetry, and second we analyze a purely off-diagonal coupling, as in the case of an O(2) family symmetry.[5] For a general family symmetry, one expects familons with both diagonal and off-diagonal couplings; a generalization to such cases is straightforward. In this subsection, we also neglect possible mismatches between the flavor and mass eigenstates, and assume that the relative angles relating the two are small. Such mismatches will be discussed in Sec. IV C. Finally, we assume that m_{ν_e}, m_{ν_μ} are negligible compared to m_{ν_τ}, as suggested from laboratory constraints as well as the corresponding masses of the charged leptons.

1. Familon with diagonal coupling

Here we consider a purely diagonal familon coupling to ν_τ, such as in models with a U(1) family symmetry acting on the third-generation lepton doublet (ν_τ, τ_L). The relevant interaction is given by

$$\mathcal{L}_f = \frac{1}{F} g_L^{\nu_\tau \nu_\tau} \partial_\mu f \overline{\nu_\tau} \gamma^\mu P_L \nu_\tau . \tag{4.1}$$

Let us first consider the case where the familon can freely escape the core of the supernova. Based on the interaction given in Eq. (4.1), potentially significant processes of familon production are the neutrino scatterings $\nu_\tau \nu_\tau \to ff$ and $\nu_\tau \to \nu_\tau f$, the latter process being allowed due to background matter effects. The familon luminosities due to these processes are given in Ref. [34]:

$$Q_f(\nu_\tau \nu_\tau \to ff) \approx 8.8\times 10^{63}\ \text{erg/sec} \times \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right)^2 \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right)^4, \tag{4.2}$$

$$Q_f(\nu_\tau \to \nu_\tau f) \approx \begin{cases} 1.6\times 10^{54}\ \text{erg/sec} \times \left(\frac{\text{MeV}}{m_{\nu_\tau}}\right)^4 \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right)^2, & m_{\nu_\tau} \geqslant 95\ \text{keV}, \\ 1.9\times 10^{60}\ \text{erg/sec} \times \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right)^2 \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right)^2, & m_{\nu_\tau} \leqslant 95\ \text{keV}. \end{cases} \tag{4.3}$$

If *either one* of the above luminosities is larger than $\approx 10^{53}$ erg/sec, the cooling process of the supernova may be dominated by familon emission, and the duration time of the neutrino pulse becomes shorter than $O(1\ \text{sec})$. Imposing the constraint that the familon luminosities given in Eqs. (4.2) and (4.3) are smaller than 3×10^{53} erg/sec, we obtain the following constraints:

$$\nu_\tau \nu_\tau \to ff: \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right) \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right)^2 \leqslant 5.8\times 10^{-6}, \tag{4.4}$$

$$\nu_\tau \to \nu_\tau f: \begin{cases} \left(\frac{\text{MeV}}{m_{\nu_\tau}}\right)^2 \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right) \leqslant 0.43, & m_{\nu_\tau} \geqslant 95\ \text{keV} \\ \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right) \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right) \leqslant 4.0\times 10^{-4}, & m_{\nu_\tau} \leqslant 95\ \text{keV}. \end{cases} \tag{4.5}$$

Familon volume emission is sufficiently small when both Eqs. (4.4) and (4.5) are satisfied.

On the other hand, if familon interactions are strong enough, familons effectively scatter off the neutrinos in the background and get thermalized and trapped in the supernova. Once this happens, a thermal sphere of familons is formed, just like the thermal neutrino sphere, and familons can only be emitted from the surface. The familon luminosity essentially obeys the formula of blackbody emission with the surface temperature of the familon sphere. The important point is that, once the familon is trapped, the familon luminosity *decreases* as the familon interaction becomes stronger. This can be understood in the following way: as the familon interaction becomes stronger, familons can be thermalized with a lower temperature. (Notice that the scattering rate increases for higher temperature.) The surface temperature of the familon sphere then decreases, and hence the luminosity is suppressed. Therefore, the familon luminosity can be small enough when the scale F is sufficiently low. Following Ref. [34] we find that the cooling through familon emission is sufficiently suppressed (i.e., is less than 3×10^{53} erg/sec) when *either one* of the following constraints is satisfied:

$$\begin{matrix} f\nu_\tau \to f\nu_\tau \\ ff \to \nu_\tau \nu_\tau \end{matrix} : \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right) \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right)^2 \geqslant 8.3\times 10^{-4}, \tag{4.6}$$

$$f\nu_\tau \to \nu_\tau: \begin{cases} \left(\frac{m_{\nu_\tau}}{\text{keV}}\right) \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right) \geqslant 85, & m_{\nu_\tau} \leqslant 1\ \text{keV} \\ \left(\frac{m_{\nu_\tau}}{\text{keV}}\right)^{1/2} \left(\frac{\text{GeV}}{F^L_{\nu_\tau \nu_\tau}}\right) e^{-m_{\nu_\tau}/2.4\ \text{keV}} \geqslant 50, & m_{\nu_\tau} \geqslant 1\ \text{keV}. \end{cases} \tag{4.7}$$

[5]Throughout our discussions, we assume that neutrinos are Majorana particles. Observations of supernova SN 1987A imply that Dirac neutrinos must be lighter than 3 keV [36] or heavier than 31 MeV [37]. As we will discuss in Sec. VI, most of the interesting mass range from a cosmological point of view is therefore excluded.

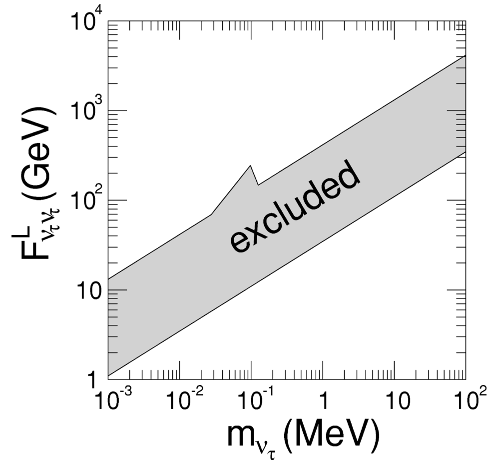

FIG. 1. Excluded region in $F^L_{\nu_\tau\nu_\tau}$ as a function of the neutrino mass m_{ν_τ}, derived from SN1987A. The bounds shown correspond to the case of diagonal familon coupling. For the off-diagonal case the bounds on $F^L_{\nu_\tau\nu_\mu}$ are very similar: the excluded region is only marginally shifted down by a factor of 1.2, and the little bump from $\nu_\tau \to \nu_\tau f$ is absent.

Of course, if the familon has strong interactions with other light particles (the photon, electron, or neutron), familons may be trapped by other processes as well. This gives additional regimes where the earlier constraints of Eqs. (4.4) and (4.5) can be evaded.

In Fig. 1, we show the upper and lower bounds on the diagonal coupling $F^L_{\nu_\tau\nu_\tau}$ as a function of the neutrino mass m_{ν_τ}. The region above the upper line is allowed because familon emission is sufficiently suppressed by the flavor scale F. This line is basically determined by Eq. (4.4); the slight bump is due to Eq. (4.5). As we can see, the lower bound on F is at most 1 TeV for the maximum allowed value of the tau neutrino mass (18.2 MeV), and it becomes less stringent as the mass becomes smaller. The lower boundary is determined by Eq. (4.6), which supercedes Eq. (4.7). The region below this line is also allowed because the familon is trapped in the core and the contribution to the cooling is again sufficiently small.

2. Familon with off-diagonal coupling

Here we discuss SN 1987A constraints on a familon which has only an off-diagonal coupling, such as in the case of an O(2) family symmetry. The low-energy Lagrangian of the model can be written as

$$\mathcal{L}_f = \frac{i}{F} g_L^{\nu_\tau\nu_\mu} \partial_\mu f (\overline{\nu_\tau} \gamma^\mu P_L \nu_\mu - \overline{\nu_\mu} \gamma^\mu P_L \nu_\tau). \tag{4.8}$$

The inclusion of ν_e in the discussion is straightforward. The familons are produced by $\nu_\tau\nu_\tau \to ff$ via t-channel ν_μ exchange, $\nu_\mu\nu_\mu \to ff$ via t-channel ν_τ exchange, or the decays $\nu_\tau \to \nu_\mu f$. Following Ref. [34] again,

$$Q_f(\nu_\tau\nu_\tau \to ff) \approx 2.2\times 10^{63}\ \text{erg/sec} \times \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right)^2 \left(\frac{\text{GeV}}{F^L_{\nu_\tau\nu_\mu}}\right)^4, \tag{4.9}$$

$$Q_f(\nu_\mu\nu_\mu \to ff) \approx 2.2\times 10^{63}\ \text{erg/sec} \times \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right)^2 \left(\frac{\text{GeV}}{F^L_{\nu_\tau\nu_\mu}}\right)^4, \tag{4.10}$$

$$Q_f(\nu_\tau \to \nu_\mu f) \approx 3.0\times 10^{62}\ \text{erg/sec} \times \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right)^4 \left(\frac{\text{GeV}}{F^L_{\nu_\tau\nu_\mu}}\right)^2. \tag{4.11}$$

We require that *all* of these familon luminosities are smaller than 3×10^{53} erg/sec, and obtain the following constraints:

$$\nu_\tau\nu_\tau \to ff,\ \ \nu_\mu\nu_\mu \to ff: \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right)\left(\frac{\text{GeV}}{F^L_{\nu_\tau\nu_\mu}}\right)^2 \leqslant 8.3\times 10^{-6}, \tag{4.12}$$

$$\nu_\tau \to \nu_\mu f: \left(\frac{\tau_{\nu_\tau}}{\text{sec}}\right)\left(\frac{\text{MeV}}{m_{\nu_\tau}}\right) \geqslant 3.3\times 10^{-5}, \tag{4.13}$$

where the lifetime of ν_τ is given by

$$\tau_{\nu_\tau}^{-1} = \frac{1}{16\pi} \frac{m^3_{\nu_\tau}}{F^{L2}_{\nu_\tau\nu_\mu}}. \tag{4.14}$$

Increasing the interaction strength further into the excluded region, familons eventually become trapped and rendered harmless again. This occurs when *any one* of the following constraints are satisfied:

$$\begin{matrix} f\nu_\mu \to f\nu_\mu,\ \ f\nu_\tau \to f\nu_\tau, \\ ff \to \nu_\mu\nu_\mu,\ \ ff \to \nu_\tau\nu_\tau \end{matrix} : \left(\frac{m_{\nu_\tau}}{\text{MeV}}\right)\left(\frac{\text{GeV}}{F^L_{\nu_\tau\nu_\mu}}\right)^2 \geqslant 1.2\times 10^{-3}, \tag{4.15}$$

$$f\nu_\mu \to \nu_\tau: (\tau_{\nu_\tau}/\text{sec})\left(\frac{\text{MeV}}{m_{\nu_\tau}}\right) \leqslant 10^{-6}. \tag{4.16}$$

As mentioned before, if the familon has strong interactions with other light particles, these interactions may lead to thermalization of familons as well, resulting in additional allowed regions for low flavor scales.

The resulting excluded region is fairly similar to that of the diagonal coupling case. The constraint from the decay process Eq. (4.13) is important only for smaller $F^L_{\nu_\tau\nu_\mu}$ or larger masses m_{ν_τ}, values that are outside our range of interest. In addition, the small bump in Fig. 1 now disappears due to the absence of the $\nu_\tau \to \nu_\tau f$ process. The dominant constraints are therefore from Eqs. (4.12) and (4.15) in the off-diagonal case, which differ from the dominant constraints of Eqs. (4.4) and (4.6) in the diagonal coupling case only by a small constant factor. The boundary of the excluded region is therefore given by the lines of Fig. 1 shifted downwards by a factor of 1.2 in F.

B. Bounds from loop-induced couplings

In the previous subsection, we considered astrophysical constraints on tree-level familon couplings to ν_τ. In addition, however, astrophysical bounds may also be used to constrain familon couplings to all other particles, as these cou-

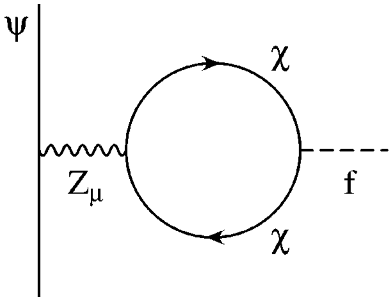

FIG. 2. The Z-f mixing diagram.

plings may induce couplings of familons to electrons and nucleons at the loop level. While these induced couplings are suppressed by the usual loop factors, the bounds on familon couplings to first generation particles are so stringent that these constraints may be strong in certain cases. In fact, we will see below that the contributions to induced couplings are proportional to fermion masses, and so these constraints are particularly relevant for couplings of familons to third generation fermions. In this subsection, we will estimate the induced couplings for various choices of the family symmetry group and determine what lower bounds on flavor scales F result from current astrophysical constraints. For simplicity, we will limit our discussion here to familons with flavor-diagonal couplings to the third generation, and ignore possible rotations relating the flavor and mass eigenstates. Extensions of this analysis to more general cases are straightforward.

To evaluate the strength of the induced coupling, we will begin by considering the low energy effective theory below the flavor scale F. In this approach, the theory is specified by the flavor symmetry, that is, the low energy derivative couplings of the familon, and no further knowledge of the mechanisms of flavor symmetry breaking is required. With the assumptions given above, the dominant contribution to the induced couplings is from the Z-f mixing graph shown in Fig. 2. Here χ is any one of the third generation particles directly coupled to the familon, and $\psi = e,\ u$, or d. (There are also additional contributions from penguin-like W diagrams, but these are suppressed by mixing angles, e.g., V_{td}^2 in the case of $\psi = d$.) Let us define the f-χ coupling as

$$\frac{1}{F}\partial_\mu f \bar{\chi}\gamma^\mu (g_L P_L + g_R P_R)\chi, \tag{4.17}$$

and the Z-χ coupling as

$$Z_\mu \bar{\chi}\gamma^\mu (g_L^Z P_L + g_R^Z P_R)\chi, \tag{4.18}$$

where $g_L^Z = g_Z(I_\chi - Q_\chi \sin^2\theta_W)$, $g_R^Z = -g_Z Q_\chi \sin^2\theta_W$, and $g_Z = e/(\sin\theta_W \cos\theta_W)$. The induced Z-f mixing from the fermion loop is divergent, and the logarithmically-enhanced contribution is

$$\begin{aligned}\mathcal{L}_{Z-f} &= \frac{2N_c}{(4\pi)^2 F}\ln\frac{\Lambda^2}{m_\chi^2}\Big[m_\chi^2 (g_L - g_R)(g_L^Z - g_R^Z) g^{\mu\nu} - \frac{1}{3}(g_L g_L^Z \\ &\quad + g_R g_R^Z)(p^2 g^{\mu\nu} - p^\mu p^\nu)\Big] Z_\mu \partial_\nu f \\ &= \frac{2N_c}{(4\pi)^2 F} m_\chi^2 \ln\frac{\Lambda^2}{m_\chi^2}(g_L - g_R)(g_L^Z - g_R^Z) Z^\mu \partial_\mu f,\end{aligned} \tag{4.19}$$

where N_c is the number of colors of the fermion χ, and Λ is the effective ultraviolet cutoff of the order of the flavor scale F. Note that $g_L^Z - g_R^Z = g_Z I_\chi$ and the fermion charge Q_χ drops out: the Ward–Takahashi identity guarantees that the $Q_\chi \sin^2\theta_W$ piece in the Z vertex gives a completely transverse vacuum polarization amplitude proportional to $p^2 g^{\mu\nu} - p^\mu p^\nu$, which vanishes when contracted with $\partial_\mu f$. The leading contributions to the induced couplings are determined by the amount of current non-conservation, i.e., the masses of the particles in the loop, and the third generation couplings therefore give the most important contributions.[6]

The mixing of Eq. (4.19) is logarithmically-enhanced and so typically gives the leading contribution if present. However, there are cases in which this term is not present or is highly suppressed. First, it may be that the amount of current non-conservation is itself suppressed by inverse powers of the flavor scale. For instance, in the singlet Majoron model [38], lepton number conservation in the low-energy theory is violated only by neutrino masses which are of order $1/F$ due to the seesaw mechanism. The neutrino loop contribution to the Z-Majoron mixing is then $1/F^3$ and highly suppressed. Second, if the familon coupling is vector-like so that $g_L = g_R$, the logarithmically-enhanced term is absent. For example, a familon coupled to the (possibly generation-dependent) baryon number current has this property. Finally, this mixing is also absent if the familon has only flavor off-diagonal couplings. In all of these cases, the contribution of Eq. (4.19) is absent or suppressed, and the leading contributions to Z-f mixing come from non-logarithmically-enhanced threshold corrections which may be of order $1/F$. Such corrections are sensitive to physics at the flavor symmetry breaking scale, and are therefore model-dependent.

With these caveats in mind, we now assume that the leading contribution given in Eq. (4.19) is present, and determine bounds on F for various flavor symmetries. The logarithmically-enhanced mixing induces an effective coupling

$$\begin{aligned}\mathcal{L}_{\text{eff}} &= \frac{2N_c}{(4\pi)^2 F} g_Z^2 I_\chi (g_L - g_R)\,\frac{m_\chi^2}{m_Z^2}\ln\frac{\Lambda^2}{m_\chi^2}\partial_\mu f \bar{\psi}\gamma^\mu (I_\psi P_L \\ &\quad - Q_\psi \sin^2\theta_W)\psi \\ &= i\,\frac{2N_c}{(4\pi)^2 F} g_Z^2 I_\chi (g_L - g_R)\,\frac{m_\chi^2}{m_Z^2}\ln\frac{\Lambda^2}{m_\chi^2} I_\psi m_\psi f \bar{\psi}\gamma_5 \psi,\end{aligned} \tag{4.20}$$

[6]Note that the radiatively-induced mixing operator can be written in the manifestly gauge invariant form $i(H^\dagger D_\mu H - D_\mu H^\dagger H)\partial^\mu f$. This operator may be present at tree level if the familon couples to ''Higgs number,'' but we will not consider this case.

where in the last step we have integrated by parts and substituted the equations of motion for ψ. For example, for a familon coupled to t_R, $N_c=3$, $I_\chi=1/2$, and $g_L-g_R=-1$. For a familon coupled to Q_L, the contributions from both t_L and b_L must be summed.

The effective coupling of Eq. (4.20) is constrained from various sources. For the case $\psi=e$, a stringent constraint is provided by red giants. Familon-electron couplings lead to additional sources of red giant cooling, which, if too large, would destroy the agreement between the observed population of red giants in globular clusters and stellar evolution theory. Such constraints have been studied extensively in the literature [39,40]. The current best upper limit on the coupling is [40]

$$g<2.5\times10^{-13} \tag{4.21}$$

for $\mathcal{L}_{\rm eff}=-igf\bar{e}\gamma_5 e$.[7] The strongest bound on the family symmetry breaking scale is for familon couplings that are dominantly proportional to m_t^2, as, for example, when a familon is coupled only to t_R. Such a case results in the bound

$$F_{tt}^R>1.2\times10^9\ \ \text{GeV}, \tag{4.22}$$

where we have taken $\Lambda=F$. Weaker, but still significant constraints are obtained if the familon coupling is dominated by m_b, as when the familon couples only to b_R. The bound in this case is

$$F_{bb}^R>6.1\times10^5\ \ \text{GeV}. \tag{4.23}$$

Notice that, in the case where the familon couples to t_R and Q_L with the same charge, the bound of Eq. (4.23) also holds for the corresponding flavor scales. However, possibly stronger bounds may also be possible if model-dependent non-logarithmically enhanced terms proportional to m_t^2 are present. If the familon contribution to the induced coupling is dominantly proportional to m_τ^2, we find the constraint

$$F_{\tau\tau}^R>2.5\times10^4\ \ \text{GeV}. \tag{4.24}$$

Similar bounds may be obtained from induced familon couplings to nucleons using constraints from supernova 1987A by rescaling the bounds on axion couplings. These constraints are somewhat more ambiguous because of the loss of coherence in axion emission due to nucleon spin fluctuations caused by scattering effects in the supernova core [41]. More realistic estimates were addressed in Refs. [42, 43]. The constraints yield results comparable to the red giant bound on electron couplings, but with larger error bars.

Finally, we stress again that these bounds are for specific flavor symmetries. For certain examples mentioned above, the logarithmically-enhanced contribution to the induced coupling is absent. For such flavor symmetries, the threshold corrections at the flavor scale must be studied separately for each model, and the flavor scale can be constrained only after the model-dependent coefficients are known. However, from the numerical estimates above, it is clear that the induced loop-level bounds can provide interesting constraints on flavor-diagonal familon couplings. Such bounds are particularly interesting for couplings to the third generation, as they are enhanced for large fermion masses. Note also that couplings to the top quark are stringently bounded and are extremely difficult to bound by other means.

C. Bounds from effects induced by flavor-mixing

In this section, we have so far parametrized and constrained possible familon couplings individually by introducing effective flavor scales, neglecting possible mismatches between flavor and mass eigenstates. However, as noted in Sec. II, when bounds on a particular familon coupling are very stringent, such as in the case of bounds on familon couplings to the first generation from supernovae, one can also obtain interesting bounds on other familon couplings from flavor-mixing effects. In this subsection, we will consider such bounds in the quark sector. In extensions of the standard model with massive neutrinos, similar arguments hold in the leptonic sector. Additional constraints may also be obtained if the gauge symmetry is enlarged.

Let us assume that the flavor and mass eigenstates coincide in the down sector. A generic familon coupling term for left-handed down-type quarks is then

$$\mathcal{L}_f=\frac{1}{F_{IJ}^L}\partial_\mu f\bar{d}_I\gamma^\mu P_L d_J, \tag{4.25}$$

where I and J are generational indices. Notice that, in this section, the familon f is a real scalar for the diagonal couplings $(I=J)$, and a complex scalar for the off-diagonal ones $(I\neq J)$. (Thus, in the off-diagonal case, there is also a Hermitian conjugate term in the Lagrangian.)

The up-type coupling required by gauge invariance in terms of down quark mass eigenstates is

$$\mathcal{L}_f=\frac{1}{F_{IJ}^L}\partial_\mu f V_{iI}^*\bar{u}_i\gamma^\mu P_L V_{jJ}u_j\equiv\frac{1}{2F_{IJ}^L}(x_u\partial_\mu f\bar{u}\gamma_\mu\gamma_5 u+\cdots), \tag{4.26}$$

where $x_u=-V_{uI}^*V_{uJ}$. Therefore, a constraint on F_{uu}^L obtained from supernova cooling through familon–nucleon coupling implies similar constraints on the expressions $F_{bs}^L/(V_{ub}^*V_{us})$, $F_{bd}^L/(V_{ub}^*V_{ud})$, and so on. Of course, different contributions to the same effective coupling $F_{uu}^{L\ -1}$ may also have opposite signs, which must be checked in the specific model under investigation.

To derive constraints on the flavor symmetry breaking scales, we must convert the quark level couplings to the effective nucleon-familon couplings of the form $\mathcal{L}_{\rm int}\sim ig_{fNN}f\bar{N}\gamma_5 N$. This can be done through a generalized Goldberger-Treiman relation. With the interaction given in Eq. (4.26), we obtain

[7] For a larger coupling, the familon may be trapped in red giants and not contribute to their cooling. Still, they can be emitted from the Sun and change its dynamics significantly. For yet larger couplings, familons may be trapped in the Sun as well, but then they contribute to the thermal transport. Combination of these constraints exclude all couplings to electrons larger than this one. See Ref. [35] for further details.

$$g_{fNN}=\frac{m_N}{F_{IJ}^L}x_u\Delta_u^{(N)}, \tag{4.27}$$

where $m_N \simeq 0.94$ GeV is the nucleon mass, and the coefficients $\Delta_q^{(N)}$ are given by [43]

$$\Delta_u^{(p)} \simeq \Delta_d^{(n)} \simeq 0.80,$$

$$\Delta_u^{(n)} \simeq \Delta_d^{(p)} \simeq -0.46,$$

$$\Delta_s^{(p)} \simeq \Delta_s^{(n)} \simeq -0.12.$$

Here, we have assumed that the flavor symmetry is anomaly-free for $SU(3)_C$. If the flavor symmetry is anomalous under $SU(3)_C$, there are anomaly-induced contributions to Eq. (4.27); see Refs. [35,43,44] for discussions of constraints on axions.

The effective couplings of Eq. (4.27) are constrained by supernova SN 1987A. In Ref. [43], the upper bound on g_{fpp} is given as a function of g_{fnn}/g_{fpp}. For simplicity, we adopt the most conservative constraint on g_{fpp},

$$g_{fpp} \lesssim 3\times 10^{-10}, \tag{4.28}$$

and use this to estimate bounds on the flavor symmetry breaking scales.

Under the assumption that only one down-type familon coupling exists at a time, that is, ignoring possible cancellations between two different contributions, we find bounds on third-generation couplings

$$F_{bb}^L > 3\times 10^4 \text{ GeV}\left(\frac{|V_{ub}|}{3.5\times 10^{-3}}\right)^2\left(\frac{\Delta_u^{(p)}}{0.80}\right), \tag{4.29}$$

$$F_{bs}^L > 3\times 10^6 \text{ GeV}\left(\frac{|V_{ub}|}{3.5\times 10^{-3}}\right)\left(\frac{|V_{us}|}{0.22}\right)\left(\frac{\Delta_u^{(p)}}{0.80}\right), \tag{4.30}$$

$$F_{bd}^L > 1\times 10^7 \text{ GeV}\left(\frac{|V_{ub}|}{3.5\times 10^{-3}}\right)\left(\frac{|V_{ud}|}{0.98}\right)\left(\frac{\Delta_u^{(p)}}{0.80}\right). \tag{4.31}$$

An even stronger constraint is obtained for F_{sd}^L:

$$F_{sd}^L > 8\times 10^8 \text{ GeV}\left(\frac{|V_{us}|}{0.22}\right)\left(\frac{|V_{ud}|}{0.98}\right)\left(\frac{\Delta_u^{(p)}}{0.80}\right). \tag{4.32}$$

This constraint is, however, weaker than the laboratory bound.

Note that the above bounds are obtained under the assumption that the mass and flavor eigenstates are identical in the down sector. If, on the other hand, these eigenstates were assumed to be identical in the up sector, familon coupling to the s- and d-quark arises due to the mixing effect, and the supernova constraints on F_{dd}^L and F_{ss}^L could be used instead. In particular, an interesting bound is derived for F_{tt}^L:

$$F_{tt}^L > 7\times 10^5 \text{ GeV}\left(\frac{|V_{ts}^* V_{ts}\Delta_s^{(p)} + V_{td}^* V_{td}\Delta_d^{(p)}|}{(4.1\times 10^{-2})^2\times 0.12 + (5.7\times 10^{-3})^2\times 0.46}\right). \tag{4.33}$$

This bound is about one order of magnitude stronger than the bound on F_{bb}^L since, in this case, the effective familon coupling to the nucleon is dominated by the s-quark contribution, and the interaction is therefore not as highly Cabbibo-suppressed as in the F_{bb}^L case.

Of course, there is no reason why the familon coupling is diagonalized in one sector. However, as the up and down sectors cannot be diagonalized in the same basis, one generally expects similar mixing-induced constraints in all cases.

V. PROSPECTS FOR FUTURE PROBES: B FACTORIES

In this section, we estimate what constraints on familon couplings may be obtained from the current CLEO data set and the upcoming CLEO III, BABAR and BELLE experiments. We make no attempt to conduct detailed experimental studies appropriate to each of these experimental settings. Rather, our intent here is to describe a number of analyses that are likely to significantly improve the present limits on familon energy scales, and, we believe, merit further study. We will begin with investigations of hadronic couplings of b quarks to familons, and then consider decays of the τ lepton to familons.

A. Bounds from B decays

In all of the experiments mentioned above, one may search for the exclusive decays $B^\pm \to (\pi^\pm, K^\pm) f$ and $B^0 \to K_s f$. These exclusive modes have smaller branching fractions than the inclusive modes $b \to (d,s) f$, but have clear experimental signatures due to their simple two-body kinematics. The form factor of

$$\langle K^-(p')|\bar{s}\gamma^\mu b(q^2=0)|B^-(p)\rangle = F_1(0)(p+p')^\mu, \tag{5.1}$$

which is necessary to calculate branching fractions, has been estimated by Colangelo *et al.* [45] to be $F_1(0)=0.25 \pm 0.03$, based on sum rules. Estimates of $F_1(0)$ based on the quark model are 0.34, 0.36, 0.30, or 0.35, depending on which quark model parameters are assumed [46]. We could not find other estimates of this particular form factor in the literature, but various estimates for $B \to \pi$ transitions give comparable but slightly larger values. This is reassuring, since they must agree in the flavor symmetric limit. The decay rate $B \to K f$ is given by

$$\Gamma(B\to Kf)=\frac{1}{16\pi}\frac{m_B^3}{F^2}g_V^2\beta^3|F_1(0)|^2, \tag{5.2}$$

where $\beta=1-m_K^2/m_B^2$. (If the coupling is purely axial, there is, of course, no contribution to $B\to Kf$; searches for decays to K^*f are required to bound such couplings.) Neglecting the mass difference between the b quark and B meson, and using the naive spectator model for the B meson decay, one finds

$$\frac{\Gamma(B\to Kf)}{\Gamma(b\to sf)}\approx|F_1(0)|^2\frac{g_V^2}{g_V^2+g_A^2}. \tag{5.3}$$

The concept of the search for such exclusive decay modes is relatively simple. After applying the standard cuts to suppress continuum $q\bar{q}$ and lepton pair events, one looks for events at the $\Upsilon(4S)$ resonance that have either an isolated K_s, or an isolated charged meson $\pi^\pm, K^\pm$ together with large missing energy. In the center-of-momentum frame, the energy of the meson $P=\pi^\pm$, $K^\pm$, or K_s must be in the narrow range

$$\frac{\sqrt{s}}{4}\left[\left(1+\frac{m_P^2}{m_B^2}\right)-\beta\left(1-\frac{m_P^2}{m_B^2}\right)\right] < E_P < \frac{\sqrt{s}}{4}\left[\left(1+\frac{m_P^2}{m_B^2}\right)+\beta\left(1-\frac{m_P^2}{m_B^2}\right)\right], \tag{5.4}$$

where $\beta=\sqrt{1-4m_B^2/s}=0.0645$, and m_P is the mass of the meson. One can also require that, after excluding the isolated energetic meson whose energy is in the above range, all the tracks and energy deposits in the calorimeters reconstruct m_B and have total energy $\sqrt{s}/2$ in the center-of-momentum frame.

Of the existing analyses, the one most similar to that described above is a search for $B^\pm\to l^\pm\nu_l$ by the CLEO Collaboration [47]. The reported upper bounds on the branching fractions are $1.5\times10^{-5}(e)$, $2.1\times10^{-5}(\mu)$, and $2.2\times10^{-3}(\tau)$. The reach for the Pf mode is expected to be worse than for $e\nu$ or $\mu\nu$, as the continuum backgrounds are larger and the detection efficiencies are worse for mesons. However, we expect the sensitivity to such meson decays to be greater than to the $\tau\nu$ mode, because the mesons have more-or-less fixed energy, unlike in the $\tau\nu$ case. π/K separation is probably difficult with the current CLEO data set [48], but this analysis may still give us an upper bound on $B(B\to\pi f)+B(B\to Kf)$ somewhere at the 10^{-4} to 10^{-3} level [49,50]. Such a constraint would be competitive with the upper bound inferred from the ALEPH $b\to s\nu\bar{\nu}$ study discussed in Sec. III C.

Particle identification will be much better at CLEO III, which will allow π/K separation, and will be even better at BABAR and BELLE. The higher luminosity at these machines will also help, and an upper bound of 10^{-5} may be possible [49]. Such a bound would imply a bound on the flavor symmetry breaking scale of $F_{bd}^V, F_{bs}^V \gtrsim 2\times10^8$ GeV.

B. Bounds from τ decays

We now turn our attention to the lepton sector. The decay rate for $\tau\to lf$ is given in Eq. (3.2). A search for $\tau\to lf$ suffers from the standard model background $\tau\to l\nu\bar{\nu}$. A conventional method for bounding the branching fraction to familons is to fit the momentum spectrum of the electron (muon) from tau decay to a linear combination of the standard model spectrum, which drops approximately linearly for large momenta, and a possible contribution from the familon mode, which is flat. The ARGUS bound quoted in Sec. III A was obtained by this method. CLEO has not reported a similar analysis. However, in a recent CLEO analysis of the Michel parameter in τ decays [51], the electron momentum distribution of 33531 $\tau^+\tau^-\to(e^\pm\nu\bar{\nu})(\pi^\mp\pi^0\nu)$ events was presented in uniform 0.25 GeV bins. This distribution may be fit beautifully by the standard model alone, and contains about 90 events in the highest momentum bin. For reference, the contribution of a familon decay mode with $B(\tau\to ef)=3\times10^{-3}$, a branching fraction near the current ARGUS limit, would contribute 28 events in each bin, leading to a significant excess at high momentum.

To see the possible sensitivity given the current CLEO data set, we generated 33531 standard model $\tau\to e+$missing events. The momentum spectrum from our Monte Carlo simulation (points with error bars), along with the predicted standard model spectrum (solid histogram), is shown in Fig. 3a. For comparison, we also plot the spectrum given a hypothetical familon branching fraction of $B(\tau\to ef)=3\times10^{-3}$ normalized to the same number of events. In Fig. 3b, we plot the ratio to the standard model prediction to make the familon contribution more visible. Note that the spectrum with the familon contribution differs considerably in the high momentum bins. By fitting the Monte Carlo data to the linear combination of the standard model and familon modes, we find that CLEO can obtain an upper bound on the familon branching fraction of 1.6×10^{-3}(95% C.L.). This would already improve the ARGUS bound [9] on $\tau\to ef$ slightly, leading to a lower bound of $F_{\tau e}>7\times10^6$ GeV. Note that in this analysis, only events with the $\pi^\pm\pi^0\nu$ decay of the other τ were used. This requirement was motivated in the original study by the desire to study spin correlations between decaying τ pairs, but is not necessary for our purpose. The statistical power of our analysis may therefore be boosted by including events with additional decay modes of the other τ.[8]

In this analysis, the systematic effects appear to be under control. The momentum dependence of the electron identification efficiency can be calibrated by using the actual data, for instance, by using radiative Bhabha events, and this calibration improves with statistics. In addition, the background is small in the above CLEO data sample. Indeed, all measurements of τ decay parameters are statistically limited and we expect a similar situation for the familon analysis. Note also that in the above analysis we simply fit to the standard model contribution, allowing its normalization to vary. In principle, one can determine this normalization by measuring

[8]Another interesting possibility is to exploit the spin correlations between decaying τ pairs by selecting events with $\omega>0$ (see Ref. [51] for the definition of ω). Such a selection enhances the right-handed τ^- (left-handed τ^+) decaying to leptons, and thereby suppresses the electron momentum spectrum at the end point. The sensitivity to familon contributions at this endpoint is then improved.

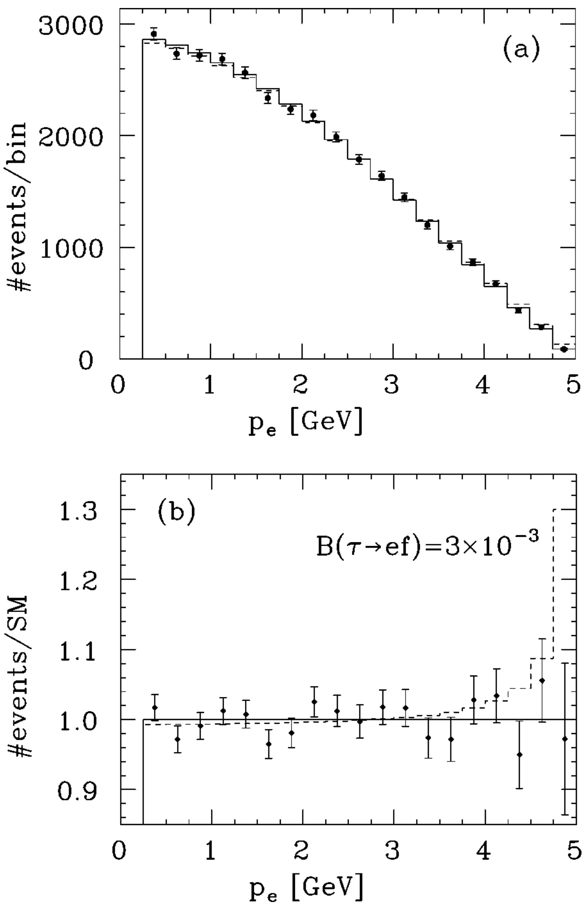

FIG. 3. (a) The e momentum spectrum from τ decays. The points with error bars are our Monte Carlo simulation for the standard model only, normalized to the size of the current CLEO event sample. The solid histogram is the standard model prediction for $\tau \to e\nu\bar{\nu}$. The dashed histogram is the predicted spectrum with $B(\tau \to ef) = 3\times10^{-3}$, again normalized to the current event sample. (b) The same as (a) but plotted as a ratio to the standard model prediction.

the efficiencies of τ identification in each decay mode through methods analogous to the multi-tag methods employed in the measurement of R_b in Z decays [52]. We therefore conclude that a dedicated analysis could well lead to an upper bound on $B(\tau \to ef)$ below the 10^{-3} level.

The $\tau \to \mu f$ mode is more difficult because the muon identification efficiency is less well-calibrated and the statistics is slightly poorer, with 21680 events in the CLEO analysis. Again, however, the uncertainties are dominated by statistical errors.[9] We therefore expect an upper bound on $B(\tau \to \mu f)$ only slightly worse than that on the ef mode.

Although we are concerned primarily with B factories in this section, we should note that the above analysis may also be applied at LEP. For example, in a recent OPAL analysis of τ polarization [53], a large sample of 25000 $\tau \to e$ events was studied. By fitting the e momentum distribution as described above, an upper bound on $B(\tau \to lf)$ at the 2×10^{-3} level could be derived. (Here, we have simply scaled the CLEO results given above to the OPAL statistics; we have checked that the momentum distributions are sufficiently similar that such an approximation is valid.) Combining the four LEP experiments, we expect an upper bound of $\sim10^{-3}$.

At the asymmetric B factories, the boosted center-of-momentum system would somewhat complicate the analysis of the electron (muon) momentum spectrum. We are not aware of any studies at these colliders. However, given that the event samples available at these machines will be much larger than the current CLEO data set, we expect BABAR and BELLE to place constraints significantly better than the 10^{-3} level, and possibly at the 10^{-4} level. The study of familon decays at these machines is extremely promising, and worthy of further study. It must be mentioned, furthermore, that the future B factories will be able to improve the upper bound on the mass of ν_τ significantly to the level of 3 MeV, or possibly even 1 MeV [54]. As we will see in the next section, the interplay between bounds on the ν_τ mass and bounds on branching ratios to familons is very interesting from a cosmological point of view.

Finally, we note that the bounds on $\tau \to lf$ branching ratios are expected to be even better at a tau-charm factory. Ref. [55] has shown that one can reach the level of $B(\tau \to ef) < 10^{-5}$ using the standard optics or even 10^{-6} using a monochromator. This would raise the lower bound on the flavor scale to $F_{\tau e} \gtrsim 10^8$ GeV. The μf mode is more difficult, and is limited by the μ/π separation capability [55]. However, a bound better than the 10^{-3} level using a RICH detector for particle identification is expected [55].

VI. IMPLICATIONS FOR NEUTRINO COSMOLOGY

Non-standard properties of neutrinos are always interesting in cosmology, and in fact, heavy unstable neutrinos are advocated in certain scenarios to obtain reasonable agreement between theory and observation. The heavy neutrino is typically taken to be the tau neutrino, and we will assume this to be the case in this section. Once a decaying neutrino is required, its decay into a lighter neutrino and a massless boson is the most harmless. Visible neutrino decays are usually severely constrained from SN 1987A. As mentioned in Sec. IV, the energy released from SN 1987A was mostly carried away in neutrinos, and the visible luminosity of SN 1987A was much smaller. However, if neutrinos decay into visible particles, such as photons or electrons, neutrinos emitted from SN 1987A that decay before reaching the earth may increase the apparent visible luminosity of SN 1987A to levels much larger than observed [56].[10] In addition, scenarios with τ neutrinos decaying into three neutrinos are also dangerous, since, in the absence of fine-tuning, such models also predict large flavor violating τ decays (such as $\tau \to 3e$) by $SU(2)_L$ gauge symmetry [13]. In particular, in the cosmological models to be described below, the resulting flavor-violating τ decay rates are already excluded by current

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[9]In the Michel parameter analysis, a large systematic uncertainty arises because the standard model muon decay parameters are not assumed. For our purposes, however, we may assume the standard model predictions and eliminate these uncertainties.

[10]Such constraints may be evaded in scenarios with sufficiently long-lived neutrinos [57].

bounds. This is because, for these three-body decays, the flavor-violating τ branching fraction is of order $(m_\tau/m_{\nu_\tau})^5\tau_{\nu_\tau}^{-1}$ [to be compared with $(m_\tau/m_{\nu_\tau})^3\tau_{\nu_\tau}^{-1}$ for the two-body familon decays], and hence is extremely enhanced for scenarios with neutrino masses in the currently allowed range.

Therefore, if we adopt a massive unstable neutrino as a solution to cosmological problems, massless bosons are good candidates for its decay products. The familon is an example of such a massless boson. (In the literature, a particular example of a familon, the Majoron, is often considered.) If neutrinos decay to familons, cosmological scenarios, which require specific ranges for neutrino masses and lifetimes, then predict rates for familon signals in future experiments, or may even be excluded from current familon bounds, assuming an absence of fine-tuning. In this section, we first review some of the potentially interesting cosmological scenarios that require massive neutrinos. Then, assuming that the massive tau neutrino decays through $\nu_\tau\to\nu_l f$, where $l=e,\mu$, we will discuss how well these scenarios may be constrained by current and future collider searches for familons.

Among the several cosmological motivations for massive neutrinos is the ‘‘crisis’’ in the standard BBN scenario. The standard BBN scenario contains only one free parameter, the baryon to photon ratio η; the abundances of the light elements are predicted once we fix η. Until a few years ago, the theoretical prediction with $\eta\sim 3\times 10^{-10}$ was in good agreement with observations. Recently, however, it has been claimed that the predictions of the standard BBN are disfavored by observations of the light element abundances [58,59]: normalizing η with the D and ^{3}He abundances, the observed ^{4}He abundance is claimed to be smaller than the standard BBN prediction. There are several arguments against this viewpoint on the observational side. For example, the apparent discrepancy vanishes if one adopts a larger systematic error in the observed ^{4}He abundance [59,60], or if the recently measured D abundance in high red-shift quasistellar object (QSO) absorber systems is regarded as a primordial one [61].[11]

On the other hand, if we regard this ‘‘crisis’’ as a genuine problem with the standard BBN theory, it can be taken as an indication of new physics beyond the standard model. There are several attempts to solve this crisis by a modification of the standard scenario [62,63,16]. Here, we concentrate on a solution that uses massive unstable neutrinos to reduce the predicted ^{4}He abundance. Since the ^{4}He abundance decreases as the energy density at the neutron freeze out time decreases, the ^{4}He abundance becomes smaller if N_ν, the ‘‘effective number of neutrino species’’ at the neutron freeze out time, is reduced. In the standard BBN, N_ν is 3, but it can be smaller if heavy neutrinos decay and effectively convert their energy density into lighter particles. For example, in the presence of the decay mode $\nu_\tau\to\nu_l f$, if all the tau neutrinos are converted into thermal familons and light neutrinos, $N_\nu\simeq 2.6$. BBN scenarios with massive neutrinos decaying to familons have been discussed in Refs. [13–16]. The most recent calculation shows that a massive neutrino with $m_{\nu_\tau}\sim 10-20$ MeV and $\tau_{\nu_\tau}\sim 10^{-2}-1$ sec can resolve the conflict between theory and observation [16].

Decaying massive neutrinos are also interesting for large scale structure formation. The standard cold dark matter (CDM) scenario, which assumes a flat universe, a scale-invariant initial spectrum, and that the universe is mostly filled with slowly moving (‘‘cold’’) particles [64], is very attractive in explaining the origin of large scale structure. However, if the normalization of the power spectrum is fixed by the anisotropy in the temperature of the cosmic background radiation observed by the Cosmic Background Explorer (COBE), the standard CDM scenario predicts too large density fluctuations at small scales ($\lambda\lesssim 100$ Mpc). Attempts to explain the scale dependence of the density perturbations include proposals of CDM with a small component of hot dark matter or with a cosmological constant [65], or scenarios with a tilted initial density fluctuation [66].

As pointed out in Refs. [18–20], CDM with late decaying neutrinos also provides a solution to this problem. If the neutrino lifetime is long enough, neutrinos dominate the energy density of the universe at the temperature $T\sim m_{\nu_\tau}$. After this stage, the mass density of the neutrino, ρ_ν, scales as T^3. Neutrinos then decay at time $t\sim(\rho_\nu^{1/2}/M_{\rm pl})^{-1}\sim\tau_{\nu_\tau}$. Once they decay, the energy density of the neutrinos is converted to radiation energy density, resulting in an increase of the radiation energy density without affecting the background photons. This then delays the time of matter-radiation equality, the matter-dominated era starts later, and, with the COBE normalization, the density perturbations at small scales are reduced. Due to the neutrino decay, the energy density of the radiation is increased by the factor $\sim\rho_\nu(T_{\rm D})/T_{\rm D}^4\sim(m_{\nu_\tau}^2\tau_{\nu_\tau})^{2/3}/M_{\rm pl}^{2/3}$, where $T_{\rm D}$ is the temperature just before the neutrino decay. As we can see, physics (approximately) depends on the combination $m_{\nu_\tau}^2\tau_{\nu_\tau}$. To obtain the correct density fluctuation at small scales, this combination must lie in the range $(m_{\nu_\tau}/\text{keV})^2(\tau_{\nu_\tau}/\text{yr})\sim 50-150$ [20].

In these scenarios, the neutrino mass must also lie in a specific interval. The neutrino must be heavier than about 50 eV; otherwise, its mass density is always smaller than (or at most comparable to) the mass density of the CDM, and the scenario does not work well. On the other hand, if the neutrino mass is above ~ 1 MeV, the neutrinos decouple from the thermal bath after becoming non-relativistic, and their number density is reduced. In this case, the constraint given above is not applicable. Furthermore, if the neutrino mass is in the range $\sim 1-10$ MeV while the lifetime is longer than ~ 1 sec, the neutrino mass density may be so large at the neutron freeze out time that ^{4}He can be overproduced. Such lifetimes and large neutrino masses are therefore also disfavored from BBN considerations [15]. In this section, we consider the mass range 50 eV$\lesssim m_{\nu_\tau}\lesssim$10 MeV, with the above caveats in mind.

To summarize, we consider the following two cosmological scenarios:

BBN: $m_{\nu_\tau}\sim 10-20$ MeV, and $\tau_{\nu_\tau}\sim 10^{-2}-1$ sec.

[11]One should note that there is another measurement of D abundances that conflicts with the one preferred by the standard BBN scenario [61]. Thus, this issue is still an open question.

Structure formation: $(m_{\nu_\tau}/\mathrm{keV})^2(\tau_{\nu_\tau}/\mathrm{yr})\sim 50-150$, with $50\ \mathrm{eV}\lesssim m_{\nu_\tau}\lesssim 10\ \mathrm{MeV}$.

These scenarios require decays to familons $\nu_\tau \to \nu_l f$. As we discussed previously, this process is related to the decay modes $\tau\to lf$ and $b\to sf$ though $SU(2)_L$ and GUT gauge symmetries, respectively. Thus, searches for these decay modes are interesting tests of these scenarios.

Let us start with the τ familon decay mode. If the decay mode $\nu_\tau\to\nu_l f$ exists, by $SU(2)_L$ symmetry, the charged τ lepton must also have flavor-changing couplings to familons:

$$\mathcal{L}_f=\frac{1}{F}\partial_\mu f(g_L^{\nu_\tau\nu_l}\overline{\nu_\tau}\gamma^\mu P_L\nu_l+g_L^{\tau l}\overline{\tau}\gamma^\mu P_L l)+\mathrm{H.c.} \quad (6.1)$$

If there is no fine-tuning, $g_L^{\tau l}\approx g_L^{\nu_\tau\nu_l}$. Notice that, if the right-handed leptons also transform under the flavor group, they also couple to familons, and such interactions may increase the rare τ decay rate. The following argument is therefore conservative. From the above Lagrangian, we obtain the decay rate

$$\Gamma(\tau\to lf)=\frac{m_\tau^3}{32\pi F_{\tau l}^{L2}}, \quad (6.2)$$

and, using Eq. (4.14), we find

$$B(\tau\to lf)=\frac{1}{2}\left(\frac{m_\tau}{m_{\nu_\tau}}\right)^3\left(\frac{\tau_{\nu_\tau}}{\tau_\tau}\right)^{-1}\left(\frac{F^L_{\nu_\tau\nu_l}}{F^L_{\tau l}}\right)^2$$
$$\simeq 8.1\times10^{-4}\times\left(\frac{1\ \mathrm{MeV}}{m_{\nu_\tau}}\right)^3\left(\frac{1\ \mathrm{sec}}{\tau_{\nu_\tau}}\right)$$
$$\times\left(\frac{F^L_{\nu_\tau\nu_l}}{F^L_{\tau l}}\right)^2. \quad (6.3)$$

In Fig. 4 contours of constant $B(\tau\to lf)$ are shown in the $(m_{\nu_\tau},\tau_{\nu_\tau})$ plane, assuming $F^L_{\nu_\tau\nu_l}=F^L_{\tau l}$.

From Eq. (6.3), we see that, in the absence of fine-tuning, the BBN scenario with decaying neutrinos predicts $B(\tau\to lf)\sim10^{-4}$ to 10^{-7}. The current bounds $B(\tau^-\to\mu^- f)<4.6\times10^{-3}$(95% C.L.) and $B(\tau^-\to e^- f)<2.6\times10^{-3}$ (95% C.L.) [9] therefore do not constrain this scenario. However, if the sensitivity of future experiments is improved by one order of magnitude or more, the predictions of this scenario may be tested, and, if the scenario is correct, exotic τ decays may be seen. (See Fig. 4.)

The scenarios motivated by structure formation are also interesting. In this case, Eq. (6.3) implies $B(\tau\to lf)\sim10^{-2}$ to 10^{-8}. Comparing this result with the current bound, part of the parameter region of this scenario is already excluded. As discussed in Sec. V B, future experiments may reach a sensitivity for $B(\tau\to lf)$ of 10^{-3} or possibly 10^{-4}. Thus, if the CDM scenario were realized, the familon decay mode is likely to be found if the neutrino is lighter than $\sim 1-10$ keV. The region disfavored by BBN [15] is also shown by a light shading.

Up to now, we have only used $SU(2)_L$ gauge symmetry to relate the neutrino-familon interaction to existing and future constraints on τ decays. However, if we assume that the same familon also couples to down-type quarks, the above cosmological scenarios may also be probed by rare b decays. Such is the case in SU(5) GUTs, where the lepton doublet and right-handed down-type quarks are in the same multiplet, and so we also have a coupling of the form

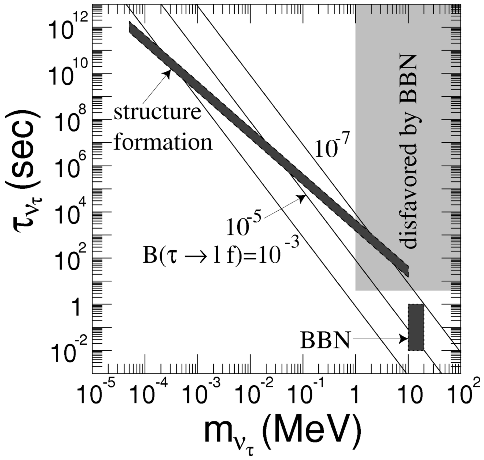

FIG. 4. Contours of constant $B(\tau\to lf)$ (10^{-3}, 10^{-5}, and 10^{-7}, from below), assuming $F^L_{\nu_\tau\nu_l}=F^L_{\tau l}$, the natural $SU(2)_L$ gauge relation in the absence of fine-tuning. The regions preferred by the BBN and structure formation scenarios discussed in the text are also shown. Note that the parameter region for structure formation with $m_{\nu_\tau}\gtrsim 1$ MeV may be unreliable because of limitations in the approximations used to derive the preferred region. The lightly shaded region is disfavored by BBN.

$$\mathcal{L}\supset\frac{1}{F}\partial_\mu f g_R^{bs}\overline{b}\gamma^\mu P_R s+\mathrm{H.c.}, \quad (6.4)$$

and similarly for d. With this Lagrangian, we obtain

$$B(b\to sf)=\frac{1}{2}\left(\frac{m_b}{m_{\nu_\tau}}\right)^3\left(\frac{\tau_{\nu_\tau}}{\tau_b}\right)^{-1}\left(\frac{F^L_{\nu_\tau\nu_l}}{F^R_{bs}}\right)^2$$
$$\simeq 7.3\times10^{-2}\times\left(\frac{1\ \mathrm{MeV}}{m_{\nu_\tau}}\right)^3\left(\frac{1\ \mathrm{sec}}{\tau_{\nu_\tau}}\right)$$
$$\times\left(\frac{F^L_{\nu_\tau\nu_l}}{F^R_{bs}}\right)^2, \quad (6.5)$$

where we have used $m_b=4.5$ GeV, and $\tau_b=\tau_B=1.6\times10^{-12}$ sec. Comparing Eq. (6.5) with Eq. (6.3), we can see that the branching ratio $B(b\to sf)$ is enhanced by about two orders of magnitude relative to $B(\tau\to lf)$. This results from an enhancement by a factor $(m_b/m_\tau)^3$, and also the fact that the total decay rate of the b quark is V_{cb} suppressed, and so is even smaller than that of the τ lepton, despite its larger mass. Contours of constant $B(b\to sf)$ are shown in Fig. 5. For the cosmologically-motivated scenarios, the ranges of $B(b\to sf)$ are 10^{-2} to 10^{-5} (BBN), and 1 to 10^{-6} (structure formation). Because the b decay rates are so enhanced, in each case, the high value in the predicted range is above bounds which can be expected from the current data. Future experiments may reach a sensitivity of $B(b\to sf)\sim10^{-4}$ [corresponding to $B(B\to Kf)\sim10^{-5}$ in Sec. V A] and will thus cover most of the preferred parameter regions.

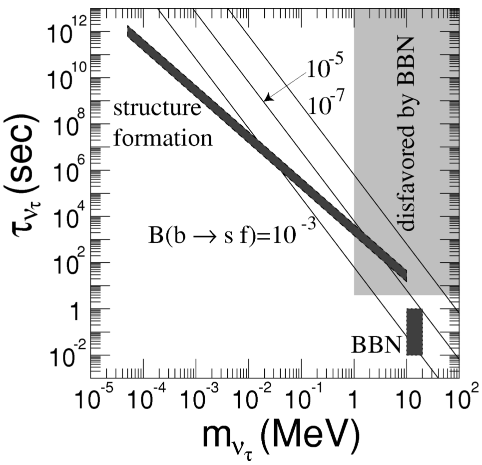

FIG. 5. As in Fig. 4, but with contours of constant $B(b\to sf)$ (10^{-3}, 10^{-5}, and 10^{-7}, from below), assuming $F^L_{\nu_\tau\nu_\mu}\simeq F^R_{bs}$, as would be the case in GUTs without fine-tuning.

Finally we comment on m_{ν_τ} measurements at future B factories. As we mentioned earlier, the upper bound on m_{ν_τ} will be significantly improved at future B factories, to the level of 1–3 MeV. The parameter region that will be covered has a significant overlap with the neutrino mass required in the above mentioned cosmological scenarios, and hence such mass measurements provide another probe of these scenarios. The BBN scenario with massive unstable neutrinos will be fully tested by the tau neutrino mass measurement in future B factories. On the other hand, for the structure formation scenario, most of the interesting parameter region ($m_{\nu_\tau}\lesssim 1$ MeV) may be covered by the search for $b\to sf$, and even if the neutrino mass is above ~ 1 MeV, this scenario can be checked by the direct measurement of the tau neutrino mass (though this region is disfavored by BBN). Therefore, in this case, measurements of the mass and the branching ratios will have complementary roles.

VII. CONCLUSIONS

If global family symmetries play a role in determining the patterns of masses and mixings of the quarks and leptons, they must be spontaneously broken. Familons (or Majorons), the massless Goldstone bosons associated with these broken symmetries, allow rare opportunities to probe the physics at very high mass scales in a multitude of low energy settings, and their discovery will signal a breakthrough in attempts to understand the flavor structure of the standard model.

The experimental investigation of familons has in the past focussed on familons coupled to the first two generations. As we reviewed, such investigations have led to stringent lower bounds on the flavor breaking scale F of $\sim 10^9$ GeV and $\sim 10^{11}$ GeV in the leptonic and hadronic sectors, respectively. In contrast, bounds for familons coupling to the third generation are much less thoroughly studied. In the lepton sector, constraints from rare τ decays lead to constraints $F \gtrsim 10^6$ GeV; in the hadronic sector, no bounds have been previously reported. The lack of study of third generation familons is conspicuous, especially in light of their cosmological relevance and the upcoming B factory experiments, which hold promise for studying b and τ decays with great precision.

Motivated by these considerations, we have presented a large and eclectic group of bounds which we believe place the most stringent constraints on third generation familon couplings. As emphasized in Sec. III, the experimental and astrophysical implications of familons vary strongly with the underlying flavor symmetry and depend on whether the flavor symmetry is real or complex, the familon couplings are axial or vector-like, and whether they are flavor-diagonal or non-diagonal. For instance, bounds on K decay in the presence of mixing effects may in general impose bounds of order 10^9 GeV on third-generation flavor breaking scales, but there are classes of models in which such bounds do not apply. It is therefore important to consider a wide variety of experimental signatures. Different signatures are also related through some well-motivated theoretical considerations: in the absence of fine-tuning, the familon couplings of particles related by gauge symmetry are expected to be similar in strength. Probes of τ decays to familons are therefore indirect probes of ν_τ familon couplings, and in SU(5) grand unified theories, these are both related to b decays.

We began by considering bounds from currently available accelerator data. Present values for neutral meson mass splittings imply bounds on the flavor scale of $\sim 10^5$ to 10^6 GeV for real familons. Bounds from neutral meson decays to leptons are significantly weaker, at the level of 10^3 GeV, and require both hadronic and leptonic couplings. More promising are bounds from exotic b decays at LEP. By extrapolating from current bounds on $b\to s\nu\bar{\nu}$, we estimate that an analysis of currently available LEP data may provide a sensitivity to $B(b\to sf)$ at the level of 1.8×10^{-3}, leading to probes of flavor scales of the order of 10^7 GeV.

Familons also have astrophysical implications, as they may lead to anomalously fast cooling of supernovae, red giants, and white dwarfs. Bounds from direct couplings to ν_τ are generally weak. However, couplings of familons to particles of the third generation may also induce couplings to electrons and nucleons radiatively or through flavor mixing effects. Bounds on such couplings are model-dependent, but may be stringent; in the simple case where a familon couples diagonally to t quarks, a bound of $F>10^9$ GeV from radiatively induced couplings may be set.

Finally, having evaluated a host of new bounds, we considered the prospects for analyses at future B factories. Such colliders are ideal experimental environments for searches for rare τ and b decay modes and are expected to have greatly improved statistics. We find that probes of branching fractions of 10^{-3} (10^{-5}) for τ (b) decays may be possible. As discussed in Sec. VI, under the assumption that the flavor scales for these decays are naturally related to the scales for ν_τ couplings, such precise probes are sensitive to parameter regions favored by various BBN and structure formation scenarios, where a massive unstable neutrino is motivated by possible discrepancies in cosmological data. In fact, parts of the parameter regions in such scenarios are already excluded, and future searches will be able to explore large portions of the cosmologically-favored parameter space. Given the present lack of analyses studying third generation familon couplings, studies at all of these experiments, and particularly the B factories, are strongly encouraged.

Note added. After the completion of this work, Ref. [67] appeared, in which the impact of decays $\nu_\tau \to \nu_e f$ on BBN was examined. Tau neutrino masses of 0.1 to 1 MeV and lifetimes of 50 to 2×10^4 s were found to be allowed, implying rates for $\tau \to e f$ and $b \to d f$ within reach of future experiments. (See Figs. 4 and 5.)

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