Probing the Origins of Neutrino Mass with Supernova Data

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We study type II supernova signatures of neutrino mass generation via symmetry breaking at a scale in the range from keV to MeV. The scalar responsible for symmetry breaking is thermalized in the supernova core and restores the symmetry. The neutrinos from scalar decays have about half the average energy of thermal neutrinos. The Bose-Einstein distribution of the scalars can be established with a megaton water Čerenkov detector. The discovery of the bimodal neutrino flux is, however, well within the reach of the Super-Kamiokande detector, without a detailed knowledge of the supernova parameters.

Type II supernovae provide some of the most extreme environments at the present cosmological epoch. In these explosions, the stellar gravitational binding energy $E_{SN} \simeq 3 \times 10^{53}$ erg is released almost entirely in a neutrino burst. After the initial collapse, the stellar core is characterized by nuclear densities and a temperature $T_{SN} \simeq 30$ MeV. The cooling of the core occurs over a period of seconds, through an explosive release of E_{SN} by neutrino emission.

Given such conditions, it is interesting to ask whether there may be new particles at a scale below T_{SN} that could be relevant in the first few seconds of the supernova explosion. In fact, supernova cooling has provided stringent bounds on many models containing light particles, like the axion [1]. However, such models do not generically yield distinct supernova signatures.

In this Letter, we study possible supernova signatures of neutrino mass generation from low energy symmetry breaking. The basic idea is that, in a typical supernova with temperature T_{SN}, a light scalar, whose vacuum expectation value (VEV) generates $m_\nu \neq 0$, may be in thermal equilibrium during the explosion. We show that this can typically be the case, providing robust signatures of this physics, which may otherwise be hard to access.

Our study is based on recent low-cutoff-scale models for neutrino mass generation proposed in Ref. [2]. In contrast to the seesaw mechanism with scales of order 10^{14} GeV, these models can be cut off at 10 TeV $\lesssim \Lambda \lesssim$ 1000 TeV, avoiding extrapolations far above the experimental frontier at around 1 TeV. Both Majorana and Dirac masses can be generated in these theories. To afford a low-cutoff-scale, a discrete gauged (anomaly-free) $Z_3^\ell \times Z_9^q$ symmetry is assumed to protect Baryon (B) and Lepton (L) numbers. Small values of $m_\nu \neq 0$ are generated when certain symmetries are spontaneously broken at or below 10 MeV, by a scalar VEV [3]. We generically denote such a scalar by φ.

A consequence of this scenario is that a Bose-Einstein gas of scalars with $T \simeq T_{SN}$ will be present during the cooling of the supernova. Thermal scalars produced in the bulk are either confined within the neutrino sphere or they decay into neutrinos, which are also confined. The φ's escaping at the surface quickly decay into neutrinos. The average energy of these neutrinos is half that of the thermal scalars. In addition, these neutrinos inherit the Bose-Einstein distribution of the parent scalars. These are distinct signatures of the presence of a light bosonic degree of freedom that couples to neutrinos. Detecting a flux of supernova neutrinos with about half the expected average energy and a Bose-Einstein thermal origin will strongly favor the type of models we consider. In this case, the supernova data will provide *a direct probe of the mechanism responsible for neutrino mass generation.*

Since the thermalized scalars are expected to have a self-coupling $\lambda \simeq 1$, the corrections to their potentials are of order $\lambda T^2 \gg \langle \varphi \rangle^2$, where $T \simeq T_{SN}$. Thus, the symmetry whose breaking gives $m_\nu \neq 0$ is restored in the supernova core. We then expect that $\langle \varphi \rangle = m_\nu = 0$ in the core, during the initial stages of the supernova explosion. A symmetry breaking *phase transition* will take place after the star has cooled to $T \lesssim \langle \varphi \rangle$. This will result in the emission of nonthermal neutrinos with $E_\nu \simeq \langle \varphi \rangle$. Later, we will briefly comment on the possibility of detecting such a signal.

To demonstrate these features, we adopt a neutrino mass model with low scale symmetry breaking, presented in Ref. [2]. In this model, Majorana neutrino masses consistent with flavor oscillation data are generated. Other models with more degrees of freedom that lead to Dirac masses from low scale symmetry breaking have also been considered [2]. To keep the discussion simple and emphasize the main features, we will not consider these latter models here. However, a similar analysis can be performed for the Dirac mass models of Ref. [2].

Without specific assumptions about the cutoff scale Λ, a gauged $Z_3^\ell \times Z_9^q$ protects B and L numbers at safe levels [2]. To obtain Majorana masses for neutrinos, a scalar φ with Z_3^ℓ charge +1 is introduced. Suppressing the generation indices, the dimension-6 operator

$$O_\varphi = \frac{\varphi L H L H}{\Lambda^2} \tag{1}$$

then yields the Yukawa coupling $y \varphi \nu_L \nu_L$, with $y \equiv (v/\Lambda)^2$ and $v \equiv \langle H \rangle \neq 0$; L and H are the standard model lepton doublet and the Higgs, respectively.

The coupling y is constrained by $0\nu\beta\beta + \varphi$ [7]: $y < 3 \times 10^{-5}$, implying that $\Lambda \gtrsim 30$ TeV. For $m_\nu \simeq 0.1$ eV, we then need $\langle\varphi\rangle \gtrsim 3$ keV. On the other hand, bounds on cosmological domain walls from broken discrete symmetries require $\langle\varphi\rangle \lesssim 1$ MeV, which yields $\Lambda \lesssim 10^3$ TeV. In this work, we will then consider 1 keV $\leq \langle\varphi\rangle \leq 1$ MeV. The upper bound on $\langle\varphi\rangle$ can easily be increased by $\mathcal{O}(1)$ factors if the domain wall network is frustrated or if the reheat temperature in the early Universe is below the symmetry breaking scale.

In a typical supernova, φ will be thermally produced in equilibrium and, consequently, the broken symmetry responsible for $m_\nu \neq 0$ will be restored. To see this, we first note that after the initial collapse of the star, the hot and dense gas of particles around this core traps neutrinos, forming a *neutrino sphere*. The ν confinement typically lasts over the cooling time, $\Delta t_{\text{cool}} \sim 10$ s, of the supernova. Assuming that the neutrinos are in thermal equilibrium inside the neutrino sphere, their number density is given by $n_\nu \sim T_{\text{SN}}^3$. The rate of the reaction $\nu_L \nu_L \to \varphi$ is given by $\Gamma_\varphi \simeq y^2 T_{\text{SN}}$ which, up to phase space factors, is of the same order as the rate for $\nu_L \nu_L \to \varphi\varphi\varphi$, given a φ 4-point coupling $\lambda \simeq 1$ [6]. We expect $y \gtrsim 10^{-7}$, with $m_\nu \simeq 0.1$ eV. This implies $\Gamma_\varphi \gtrsim (10^{-8}$ s$)^{-1}$.

Thus, $\Gamma_\varphi^{-1} \ll \Delta t_{\text{cool}}$. Note that Γ_φ also sets the rate for $\nu_L \varphi \to \nu_L$ and $\nu_L \varphi \to \nu_L \varphi\varphi$, mediated by the t-channel exchange of φ. The mean free path d_φ for the interactions of the φ's produced inside the neutrino sphere is then given by $d_\varphi \sim \Gamma_\varphi^{-1} \lesssim 1$ m. The size R_ν of the neutrino sphere is comparable to that of a neutron star and we have $R_\nu \simeq 50$ km $\gg d_\varphi$. Given all these considerations, we conclude that the neutrinos and the φ's come to thermal equilibrium inside the neutrino sphere, well before the star begins cooling. This means that the number density of φ's is given by $n_\varphi \sim n_\nu \sim T_{\text{SN}}^3$.

To study finite temperature effects, we write down a Z_3^ℓ-invariant potential for φ below the weak scale

$$V(\varphi) = y\varphi\nu_L\nu_L - \mu^2\varphi\varphi^\dagger + m\varphi^3 + \lambda(\varphi\varphi^\dagger)^2 + \text{h.c.} \quad (2)$$

In a typical theory, $\mu \simeq m \simeq \langle\varphi\rangle$, and $\lambda \simeq 1$. Since experimental constraints require $y < 3 \times 10^{-5}$, the Yukawa coupling is much smaller than other typical couplings of the φ system. Hence, Yukawa contributions to thermal corrections of $V(\varphi)$ are ignored in our analysis.

Once in equilibrium, the φ field obtains a thermal mass $\sim \lambda T^2$. At $T \simeq T_{\text{SN}}$, this correction is much larger than the typical negative mass squared responsible for symmetry breaking: $\lambda T_{\text{SN}}^2 \gg \mu^2$. Hence, the spontaneously broken symmetry is restored within the neutrino sphere.

The temperature of the neutrino sphere will eventually fall below the critical temperature T_c at which the symmetry is broken. Here, we note that due to the presence of the term $m\varphi^3$, the transition back to the broken phase may be first order. Quite generically, we expect $T_c \simeq \langle\varphi\rangle$. Once $T < T_c$, the field φ will roll or tunnel to its vacuum value.

In doing so, φ will oscillate about its VEV and radiate its energy in neutrinos. These neutrinos will carry a typical energy of order $m_\varphi \simeq \langle\varphi\rangle$.

To estimate the potential signal, we first note that the total energy stored in the symmetric phase E_φ is given by $\sim \langle\varphi\rangle^4 V$; V is the volume of the neutrino sphere. Comparing E_φ to E_{SN} for $\langle\varphi\rangle = 1$ MeV, we see that only about $10^{-6} E_{\text{SN}}$ will be emitted directly during the phase transition. To have a chance to detect such a tiny neutrino emission, only the largest available detector with a very low threshold can be used. The IceCube detector has excellent capabilities to study supernova neutrinos, although its single event energy threshold is around 100 GeV. Instead, supernova neutrinos will be detected only by the ice glow their charged current interactions produce, which leads to an increase in the noise rate of the photomultiplier tubes [8]. The signal in the IceCube detector is thus proportional to E_ν^3, where a factor E_ν^2 is due to the energy dependence of the cross section and an additional factor E_ν accounts for the fact that the number of Čerenkov photons is directly proportional to E_ν. The number of neutrinos emitted in the transition scales as $\langle\varphi\rangle^3$. Since $E_\nu \simeq \langle\varphi\rangle$, the total signal scales like $\langle\varphi\rangle^6$. On the other hand, the minimum signal strength needed for a significant detection depends on the noise of the phototubes f and on the time interval of the neutrino emission Δt_φ like $1/\sqrt{f\Delta t_\varphi}$. Using a realistic set of numbers to describe the IceCube detector from Ref. [8], we find that a minimal signal for the phase transition requires $\langle\varphi\rangle \gtrsim 5$ MeV and $\Delta t_\varphi \lesssim 10^{-4}$ s. This would require that the usual upper bounds on $\langle\varphi\rangle$ be somewhat relaxed and also that the whole neutrino sphere undergo the phase transition at the same time, in a *neutrino flash*.

Next, we will present our quantitative results for the expected size of the signal from φ-gas decays at the neutrino sphere. In thermal equilibrium, φ's will track the neutrino density and temperature closely; i.e., they are confined within the neutrino sphere by the $\varphi \leftrightarrow \nu\nu$ reactions. Neutrinos escape from the surface of the neutrino sphere. Similarly, φ's also escape, but decay into neutrinos promptly, as they have a lifetime $\tau \sim T/(\lambda m_\nu^2) \sim 10^{-6}$ s in the star's rest frame. Here we use T/m_φ to account for time dilation and $m_\nu = y\langle\varphi\rangle$. Thus, the neutrino sphere will radiate Fermi-Dirac thermal neutrinos, as well as those from decays of the Bose-Einstein φ gas.

The relative luminosity is then determined by the energy fractions carried by neutrinos and φ's. The two-component gas of neutrinos and φ's will have its energy equipartitioned into all available degrees of freedom: 3 fermionic and 1 bosonic. Thus, the number of effectively massless degrees of freedom is $3 \times (7/8) + 1 = 29/8$; a fraction $F_B = 8/29 \simeq 0.28$ of the total neutrino luminosity will be emitted from a Bose-distributed gas. As the most important detection channel is inverse β decay, we next ask: how large is the φ-decay contribution to the $\bar{\nu}_e$ flux?

The decay of the φ particles produces mass eigenstates instead of flavor eigenstates, hence the branching ratio into a flavor α depends on the hierarchy of masses m_i and the relevant mixing matrix elements $U_{\alpha i}$: $\mathrm{Br}(\nu_\alpha) \propto \sum_i |U_{\alpha i}|^2 m_i^2$. Here, we note that neutrino oscillation data have provided some information on $U_{\alpha i}$ [9]: $|U_{e1}|^2 \simeq 0.7$ and $|U_{e2}|^2 \simeq 0.3$ whereas for $|U_{e3}|^2$ there is only an upper bound of 0.05 at the 3σ level. Since $\mathrm{Br}(\nu_\alpha)$ is proportional to m_i^2, we can distinguish three cases. The first one is normal hierarchy; i.e., m_3 is much heavier than m_1 and m_2. In this case basically only the state m_3 is produced and $\mathrm{Br}(\bar\nu_e) = 1 \times |U_{e3}|^2 \leq 0.05$. In the case of inverted hierarchy, m_1 and m_2 have nearly the same mass and are much heavier than m_3, which gives $\mathrm{Br}(\bar\nu_e) = (1/2)|U_{e1}|^2 + (1/2)|U_{e2}|^2 \simeq 1/2$. In the case of degenerate neutrinos, all three masses are approximately the same and $\mathrm{Br}(\bar\nu_e) = (1/3)|U_{e1}|^2 + (1/3)|U_{e2}|^2 + (1/3)|U_{e3}|^2 \simeq 1/3$. The branching ratio of the φ decay determines the flavor ratio for the bosonic flux component: $\eta^B_{\nu_\alpha} = \mathrm{Br}(\nu_\alpha)$.

In a simple blackbody approximation to the neutrino emission from the proto–neutron star, there would be again equipartition and thus the flavor ratio for the Fermi-Dirac component would be $\eta^D_{\nu_\alpha} = 1/3$. Certainly, the effect in $\bar\nu_e$ would be very small for normal hierarchy in the absence of neutrino oscillations. More detailed studies of neutrino emission including various levels of microphysical detail show a more or less pronounced difference from the blackbody Ansatz and deviations in $\eta^D_{\nu_\alpha}$ from $1/3$ up to a factor of 2 seem possible [10].

In general, the $\bar\nu_e$ flux arriving at the detector $\Phi_{\bar\nu_e}$ is a sum of 6 different initial fluxes at the supernova $\phi_{\bar\nu_\alpha}$

$$\Phi_{\bar\nu_e} = N \sum_{\bar\nu_\alpha} P_{\bar\alpha\bar e}(F_D \eta^D_{\bar\nu_\alpha} \phi^D_{\bar\nu_\alpha} + F_B \eta^B_{\bar\nu_\alpha} \phi^B_{\bar\nu_\alpha}), \qquad (3)$$

where $P_{\bar\alpha\bar e}$ denotes the probability of neutrinos of flavor $\bar\alpha$ to arrive as $\bar\nu_e$ at the detector. N is a normalization factor accounting for the energy and distance of the supernova. Here, there is basically no difference between $\bar\nu_\mu$ and $\bar\nu_\tau$. Both interact with the neutron-star material only via neutral current interaction and any differences in their initial spectrum would be eliminated by neutrino oscillation via the atmospheric angle. Thus, it is sufficient to treat $\bar\nu_\mu$ and $\bar\nu_\tau$ as one flavor $\bar\nu_x$ with twice the flux. Having now an effective two flavor problem we can write $P_{\bar x \bar e}$ as $1 - P_{\bar e \bar e}$. Since $F_B + F_D = 1$, Eq. (3) yields

$$\Phi_{\bar\nu_e} = N\{P_{\bar e\bar e}[(1 - F_B)\eta^D_{\bar\nu_e}\phi^D_{\bar\nu_e} + F_B\eta^B_{\bar\nu_e}\phi^B_{\bar\nu_e}]$$
$$+ (1 - P_{\bar e\bar e})[(1 - F_B)\eta^D_{\bar\nu_x}\phi^D_{\bar\nu_x} + F_B\eta^B_{\bar\nu_x}\phi^B_{\bar\nu_x}]\}. \quad (4)$$

The η^B's are determined by the solar mixing angle and the mass ordering. Therefore, we will discuss the different possibilities for the mass ordering. However, within each case, we will assume each η^B to be known, since the solar mixing angle is well constrained. $P_{\bar e\bar e}$ also depends on the solar mixing angle, the mass ordering, and additionally on

the small angle θ_{13}. The value of $P_{\bar e\bar e}$ is determined by the adiabaticity of the neutrino evolution at the Mikheyev-Smirnov-Wolfenstein resonance which corresponds to the larger, atmospheric mass splitting. The adiabaticity itself depends on the value of θ_{13}; for a detailed derivation see, e.g., Ref. [11].

Since there is only an upper bound on θ_{13}, we will consider two extreme cases: either $\sin^2\theta_{13}$ is much smaller than 10^{-3} or much larger. In the first case, $P_{\bar e\bar e} = \cos^2\theta_{12} \simeq 0.3$, irrespective of the mass ordering; in the second case $P_{\bar e\bar e} = \cos^2\theta_{12}$ for normal mass hierarchy and $P_{\bar e\bar e} = 0$ for inverted hierarchy. Thus, at least a fraction $1 - \cos^2\theta_{12} \simeq 0.7$ of the $\bar\nu_x$ flux will contribute to the observable $\bar\nu_e$ flux. Therefore, there is always a sizable contribution to the signal by neutrinos from φ decay even for the normal hierarchy, where nearly all φ's decay into $\bar\nu_x$, since at least 70% of $\bar\nu_x$ will oscillate into $\bar\nu_e$.

The spectral distribution for the initial fluxes in the Fermi-Dirac case is given by the usual expression. However, for the neutrinos originating in φ decays, we have to consider that each φ particle will decay into two neutrinos, with flat energy distributions in the star's rest frame. Hence, the distribution functions we use are

$$\phi^D_{\bar\nu_\alpha}(E; T_{\bar\alpha}) = \frac{E^2}{e^{E/T_{\bar\alpha}} + 1} \qquad (5)$$

and

$$\phi^B_{\bar\nu_\alpha}(E; T_{\bar\alpha}) = 2 \int_E^\infty \frac{dE'}{E'}\left(\frac{E'^2}{e^{E'/T_{\bar\alpha}} - 1}\right), \qquad (6)$$

where the factor of 2 in ϕ^B accounts for 2 neutrinos per decay. We implicitly assume that the temperature of the φ gas and the Fermi-Dirac gas are the same, but we allow for different temperatures of $\bar\nu_e$ and $\bar\nu_x$. The signature for the presence of a φ gas is thus the observation of up to four different distributions which compose the signal. Two of them are the usual Fermi-Dirac distributed $\bar\nu_e$ and $\bar\nu_x$ contributions, and we have another two contributions which result from φ decay. The latter two have an average energy which is approximately only $1/2$ of that corresponding to the supernova temperature.

Our simple ansatz does not include possible changes in the spectral shape with respect to the pure Fermi-Dirac or Bose distributions. Those changes arise due to the energy dependent position of the neutrino sphere and due to diffusion processes on the way out from the neutron star. Since neutrino density is relatively low outside the neutrino sphere, Pauli blocking can be ignored [10]. Therefore, one naively expects that the Bose and Fermi-Dirac components do not interact on their way out. Thus, transport or diffusion should affect both components in a similar fashion. To understand the effects of the φ gas on the position of the neutrino sphere(s) and its energy and flavor dependence requires a detailed calculation along the lines of Ref. [10] and is beyond the scope of this Letter. However, it seems

TABLE I.   Required number of events for a 3σ discovery.

Hierarchy $P_{\bar{e}\bar{e}}$	Normal $\cos^2\theta_{12}$	Inverted $\cos^2\theta_{12}$	0
Perfect knowledge	538	152	294
N unknown	1107	678	682
N and $T_{\bar{\nu}_e}$ unknown	1157	709	682
N, $T_{\bar{\nu}_e}$, and $T_{\bar{\nu}_x}$ unknown	10 502	7397	709
N, $T_{\bar{\nu}_e}$, $T_{\bar{\nu}_x}$, and $\eta^D_{\bar{\nu}_e}$ unknown	13 449	7432	709

unlikely that such effects could conspire to destroy the four-component feature.

It remains to estimate the size of the event sample needed to detect the bimodal nature of the $\bar{\nu}_e$ spectrum. To this end, we convert the flux in Eq. (4) to an event rate by convolving it with the cross section for inverse β decay which is approximately given by $\sigma = 9.52 \times 10^{-44}E_\nu^2$ cm^2 MeV^{-2}, with a Q value of 1.8 MeV. We can now compute the event spectrum for a given set of parameters $T_{\bar{\nu}_e}$, T_{ν_x}, $\eta^D_{\bar{\nu}_e}$, $P_{\bar{e}\bar{e}}$, F_B and ask how well we can determine or constrain F_B from a fit to these simulated data. We will use the inverse β-decay cross section from Ref. [12] and assume a typical detector threshold of 7 MeV [13]. We use the supernova input parameters

$$T_{\bar{\nu}_e} = 4.8 \text{ MeV}, \qquad T_{\nu_x} = 4.8 \text{ MeV}, \qquad \eta^D_{\bar{\nu}_e} = 0.66,$$
(7)

corresponding to the simulation in Ref. [10]. With $P_{\bar{e}\bar{e}} = \cos^2\theta_{12}$, a supernova with a total energy release of 3×10^{53} erg at a distance of 10 kpc, yields $\sim 10\,000$ events, in a detector with a fiducial mass of 22.5 kt, like Super-Kamiokande. We simulate our data assuming that indeed there is a φ gas, i.e., $F_B = 0.28$.

We perform a least squares fit to the simulated data in the usual way by constructing a χ^2 function and asking how many events are needed to exclude the absence of the φ gas, $F_B = 0$, at 3σ confidence level. More technically, we want to know how large $\Delta\chi^2(F_B = 0)$ is per event. The fit is constrained to $T_{\bar{\nu}_e}, T_{\bar{\nu}_x} \geq 3.8$ MeV, which is a conservative range [10]. The number of events needed for a 3σ discovery, under various assumptions about our knowledge of the supernova, is shown in Table I. From this table, we see that the Super-Kamiokande detector is quite capable of detecting the presence of the φ gas without the need for a detailed model of the supernova physics.

Establishing the Bose-Einstein origin of the flux from φ decay will require a large data sample. Here, one can perform an analysis similar to the one above, using a Fermi-Dirac distribution for the scalar gas when testing the χ^2 difference in the fit to the simulated data. We find that 10^5–10^6 events are needed to reach a conclusion at 3σ. For this purpose, our estimate suggests, one of the megaton water Čerenkov detectors will be needed.

In this Letter, we have presented supernova signatures of neutrino mass models where new massive scalars with VEV's in the range keV–MeV generate $m_\nu \neq 0$. We have shown that the scalars will typically come to thermal equilibrium during the initial stages of supernova cooling. The subsequent decay of the scalars yields a neutrino flux with roughly half the average energy of thermally emitted neutrinos. The lower energy component of the neutrino flux encodes the Bose distribution of the parent scalars. A bimodal energy distribution for each flavor is a distinct and robust signature of the models we consider. This effect could be readily detectable at the 3σ level at the Super-Kamiokande detector, for a typical galactic supernova. Observation of these signatures may offer the only direct access to the physics of neutrino mass generation.

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