Limits to radiative neutrino decay from SN 1987A

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We calculate limits to the properties of massive, unstable neutrinos using data from γ-ray detectors on the Pioneer Venus Orbiter (PVO) Spacecraft. The absence of a γ-ray signal in the PVO detector constrains the branching ratio to photons (B_γ), mass (m_ν), and radiative lifetime ($\tau_\gamma=\tau/B_\gamma$). For low-mass ($m\lesssim T\sim 8$ MeV) neutrinos decaying $\nu\rightarrow\nu'\gamma$, $B_\gamma<3\times10^{-7}$ for $m_\nu\tau\lesssim10^6$ keVsec, and $B_\gamma<2\times10^{-13}m_\nu\tau$/keV sec for $m_\nu\tau\gtrsim10^6$ keV sec; limits for high-mass neutrinos are somewhat weaker due to Boltzmann suppression. We also calculate limits for decays that produce γ rays through the bremsstrahlung channel, $\nu\rightarrow\nu' e^+e^-\gamma$. With one exception, the PVO limits are roughly comparable to those from an analysis of data from the Solar Max Mission (SMM) Satellite (which observed at higher γ-ray energies but for a much shorter time). For neutrino mass states that are nearly degenerate, $\delta m^2/m^2\sim0.1\ll1$, our limits for the mode $\nu\rightarrow\nu'\gamma$ become more stringent by a factor as large as $m^2/\delta m^2$, because more decay photons are shifted into the PVO energy window. For this same reason, SMM cannot constrain this case. [S0556-2821(97)05912-2]

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I. INTRODUCTION

The occurrence of Supernova 1987A (SN 1987A) in the Large Magellanic Cloud has proven to be among the most fruitful experiments in the heavenly laboratory, confirming ''known'' physics and constraining new physics. Aside from its obvious impact upon the study of the late stages of stellar evolution in general and upon supernova physics in particular, models for SN 1987A have become a test bed for the study of the couplings of light particles (e.g., neutrinos, axions) to ordinary matter [1]. In this work, we discuss limits on the properties of neutrinos based upon their thermal emission from the hot nascent neutron star.

When a supernova occurs, the bulk of the binding energy of the neutron star ($\sim3\times10^{53}$ erg) is released in neutrinos, as predicted by theory and confirmed by the observation of a neutrino burst from SN 1987A. The temperature of electron neutrinos is about $T_\nu\simeq4.5$ MeV; μ and τ neutrinos are predicted to have a higher temperature, $T\approx8$ MeV, because they couple to the prevalent electrons only through neutral current interactions [2–4]. If at least one species of neutrinos is massive, unstable and couples to the photon, then some of these neutrinos will decay en route to photons, potentially detectable as MeV γ rays. At the time of the supernova burst's arrival at the Earth, there were several satellites operating in the solar system capable of detecting the decay photons in the course of their watch for γ-ray bursts. Analyses of the data from one of these detectors, on board the Solar Max Mission (SMM) Satellite, have already been presented [5,6]; here, we examine the data from the Gamma Burst Detector on the Pioneer Venus Orbiter (PVO) [7]. While the PVO detector was smaller, and its energy window (0.1–3 MeV) not well matched to that of the supernova neutrinos, it had 4π acceptance and was in an environment free of the Earth's radiation belts (leading to lower backgrounds). In addition, more high-quality data is available (>8000 sec vs 10 sec for SMM). These factors combine to give limits that are comparable to, but slightly weaker than, those from SMM [6], cf. Sec. III. For the case of nearly degenerate neutrinos (Sec. VI), most of the γ-ray flux is at a lower energy, and the PVO limits are correspondingly stronger.

The paper is organized as follows. In the next section an exact formula for the expected γ-ray flux is derived and important approximations are developed. In Sec. III, the PVO data are discussed and rigorous limits are derived in the simplest regime. The next four sections build upon these results, expanding to more complicated regimes. The final section is a brief summary.

II. γ-RAY SIGNAL

The fluence of γ rays from decaying massive neutrinos radiated from the nascent neutron star depends upon the particular decay channel. Here, we will consider the simplest two-body decay, $\nu\rightarrow\nu'\gamma$, with a low-mass daughter neutrino and a short lifetime; in Secs. IV–VII below, we allow further complications which have not been examined as extensively and for which the PVO data is better suited.

We can write the expected fluence of γ rays from decaying neutrinos with mass m_ν and mean lifetime τ as [8]

$$dN=\frac{B_\gamma L_\#(E)}{4\pi D^2}dEn(\mu)d\mu\frac{e^{-t_d/\gamma\tau}}{\gamma\tau}dt_d\delta\bigg(t-t_d\bigg[1-v\mu+\frac{D}{t_d}\{\sqrt{1-(vt_d/D)^2(1-\mu^2)}-1\}\bigg]\bigg)dt, \quad (1)$$

where B_γ is the fraction of decays that produces a γ ray. The first factor is the overall flux of neutrinos from a supernova at a distance D. $L_\#(E)$ is the differential number flux of neutrinos, so $E_T=\int dEEL_\#$ is the total luminosity in neutrinos. The second factor gives the fraction that decays into a ''laboratory-frame'' angle $\arccos\mu$. The third factor gives the fraction that decays at time t_d. Finally, the δ function selects the photons with a given t_d, E, μ that arrive at a time t after the arrival of massless neutrinos at the detector. The Lorentz factor is $\gamma=E/m_\nu$ and the speed $v=\sqrt{1-\gamma^{-2}}$.

The function $n(\mu)$ depends on the distribution of daughter photons in the neutrino rest frame and, therefore, on the particular decay channel. First, we consider the two-body decay, $\nu\rightarrow\nu'\gamma$. Because the neutrino is a spin-1/2 particle, and the photon a spin-1 particle, this reaction can proceed in one of two ways: with the helicity of the daughter neutrino parallel or antiparallel to the photon helicity. From quantum mechanics, then, the distribution of the photon in the rest frame of the parent will be proportional to either $(1\pm\overline{\mu})/2$, where $\overline{\mu}$ is the cosine of the rest-frame angle between the directions of the parent neutrino and the photon. Transforming into the laboratory frame gives the distribution $n(\mu)$. Note that we have assumed neutrinos are emitted in an instantaneous burst; as long as the actual duration is small compared to the timing resolution of the detector, which is the case, this is a good approximation.

Because we are not interested in the decay angle, but rather the photon energy, we write $n(\mu)d\mu=f(E,k)dk$, where k is the γ-ray energy, related to the decay angle by

$$\mu=\frac{\gamma}{\sqrt{\gamma^2-1}}\left(1-\frac{m}{2\gamma k}\right). \tag{2}$$

This gives

$$f(E,k)=\frac{1}{(Ev)^2}(Ev\mp E\pm 2k), \tag{3}$$

for each of the helicity possibilities. [For reference, an isotropic decay would give $f(E,k)=1/(Ev)=1/p$, where p is the neutrino momentum.] For ultrarelativistic neutrinos $(m_\nu\ll T_\nu)$, $v\approx 1$,

$$f(E,k)=\begin{cases} 2k/E^2 & \text{no flip,}\\ 2(E-k)/E^2 & \text{flip,}\\ 1/E, & \text{isotropic.}\end{cases} \tag{4}$$

Each of these should be multiplied by a Heaviside function $\Theta(E-k)$ to require that the daughter photon be less energetic than the parent. Further, we ensure that the decay does not occur inside of the progenitor envelope which would considerably alter the energetics of the explosion and lead to an independent constraint which is important for short lifetimes [9].

The factor inside the δ function in Eq. (1) is especially complicated. This is because at any given time, the detector is receiving photons from neutrinos that have decayed on a surface with a complicated shape that is approximately ellipsoidal with the supernova at one focus and the detector at another, and further complicated if the speed $v<1$ (for massive neutrinos). This includes photons that have left the supernova pointing far away from the detector but which decayed at large angles toward the detector. Obviously, for low-mass neutrinos which leave the supernova at highly relativistic speeds, the fraction that takes such a path is very small. To simplify this expression, we shall require that $t_d\ll D\sim 5\times 10^{12}$ sec, the assumption that most of the neutrinos decay well before they reach the earth. In Sec. VII we discuss long lifetimes, in which case the flux is greatly reduced and the limits are correspondingly weaker.

In the present limit, the δ function becomes simply $\delta[t-t_d(1-v\mu)]=(2\gamma k/m)\delta(t_d-2\gamma kt/m)$, and we can perform the integration over t_d:

$$dN=\frac{B_\gamma L_\#(E)}{4\pi D^2}f(E,k)\frac{2k}{m_\nu\tau}e^{-2kt/m_\nu\tau}dtdEdk. \tag{5}$$

(Similar expressions have also been derived in Refs. [6,10].) We shall assume that the neutrino-number luminosity is given by a zero-chemical-potential Fermi-Dirac (FD) distribution with known temperature and total energy, a reasonable approximation [3]. For now, we consider low-mass neutrinos (i.e., $m_\nu\ll T_\nu$) where

$$L_\#(E)=\frac{120}{7\pi^4}\frac{E_T}{T_\nu{}^4}\frac{E^2}{1+e^{E/T_\nu}}, \tag{6}$$

where $E_T\simeq 10^{53}$ erg is the total energy in one species of massless neutrinos. We treat the case $m_\nu\gtrsim T_\nu$ in Sec. V below.

Finally, we can integrate the above expression over neutrino energy E and over one time bin, from t to $t+\delta t$, to get an expression for the spectrum of photons incident on the detector during that time interval to obtain

$$\begin{aligned}\phi(k,t)&=\int_t^{t+\delta t}\frac{dN}{dkdt}dt\\ &=\frac{B_\gamma}{4\pi D^2}\frac{240}{7\pi^4}\frac{E_T}{T_\nu{}^2}h(k/T_\nu)e^{-2kt/m_\nu\tau}(1-e^{-2k\delta t/m_\nu\tau}).\end{aligned} \tag{7}$$

In this expression, the function $h(k/T_\nu)$ results from the integration over the neutrino energies. It is of order unity for the parameter ranges of interest, and it is largest in the case of ''no flip,'' which we will assume from now on since it gives the most conservative estimates of the parameters. In that case, it is given by $h(y)=y\ln(1+e^{-y})$.

Although this signal depends nonlinearly on the parameter $m_\nu\tau$, the expression simplifies when $m_\nu\tau$ is much greater than kt or much less than $k\delta t$, where k is a typical photon energy,

$$\phi(k,t)=\frac{B_\gamma}{4\pi D^2}\frac{240}{7\pi^4}\frac{E_T}{T_\nu{}^2}h(k/T_\nu)\begin{cases}\delta_{t,0}, & m_\nu\tau\ll k\delta t,\\ 2k\delta t/m_\nu\tau, & m_\nu\tau\gg kt.\end{cases} \tag{8}$$

For small $m_\nu\tau$, there is no appreciable relativistic time delay before the decay of the neutrinos, so essentially all of the daughter photons arrive in the first time bin. In the case of large $m_\nu\tau$, the flux is essentially constant over the time of the

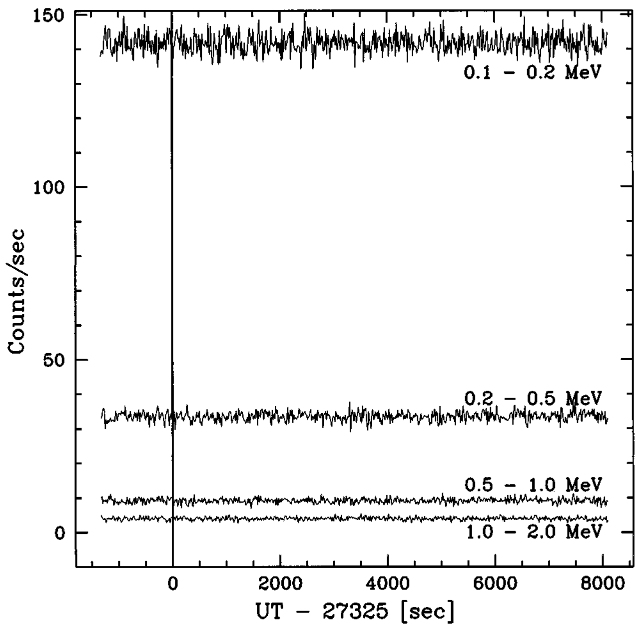

FIG. 1. The PVO GBD data for the time immediately before and after the arrival of the neutrino pulse from SN 1987A at the PVO spacecraft (UT = 27325). Time bins are either 12 or 16 sec; we show the average counts per second in each bin, for each energy channel, as marked.

observations, so the signal is proportional to the width of the time bin. Only in the latter case does the fluence actually depend on the value of $m_\nu\tau$.

In order to calculate the expected signal, we must fold the photon spectrum with the detector response function. The signal expected in the ith energy channel is

$$S_i(t)=\int dkR_i(k)\phi(k,t)=\sum_j R_{ij}\phi_j(t), \tag{9}$$

where R_{ij} is the response of detector i to energy bin k_j, and $\phi_j(t)$ is the theoretical spectrum averaged over energy bin j at time (or time bin) t.

III. γ-RAY DATA AND ANALYSIS

To obtain our limits we use data from the Pioneer Venus Orbiter Gamma Burst Detector (PVO GBD) [7,11] around the time of the supernova. The GBD has four energy channels, roughly 100–200 keV, 200–500 keV, 500–1000 keV, and 1–2 MeV. The supernova was propitiously directly overhead at the time, giving the maximum effective area. We have data for about 1500 sec prior to the arrival of the supernova neutrinos at Venus (for calculating the background), and for 8000 sec after, for time bins of either 12 or 16 sec in duration. We show the data in Fig. 1. We have verified that there is no clear signal in any of the four channels: the data is consistent with a constant Poisson rate in each detector.

To calculate limits on our parameters B_γ and $m_\nu\tau$, we use the expected signal $S_i(t)$ and our measurement of the background over a time interval t_b in each detector to construct a likelihood function given the observed γ-ray counts in each detector. We assume that $S_i(t)$ gives the mean of a Poisson process governing the detected number of counts; the rates for each detector are high enough to be well described by a normal distribution, which we use for ease of calculation (in the 1–3 MeV bin, with the lowest fluence, there are approximately 40 counts per bin). This gives a likelihood function

$$\mathcal{L}(\theta)=\prod N[D_{ij};b_i\delta t_j+S_{ij}(\theta),\sigma_{ij}^2], \tag{10}$$

where

$$N(x;\mu,\sigma^2)=\frac{1}{\sqrt{2\pi\sigma^2}}\exp\left[-\frac{1}{2}\frac{(x-\mu)^2}{\sigma^2}\right] \tag{11}$$

is the normal distribution, b_i is the background in channel i (observed for a time t_b), δt_j is the length of time bin j, and D_{ij}, S_{ij} are, respectively, the observed and theoretical signal in those time bins, calculated with the set of parameters represented by θ. Finally, the variance is given by $\sigma_{ij}^2=S_{ij}+b_i\delta t_j(1+\delta t_j/t_b)$, the sum of the theoretical variance of the signal and that due to the background rate.[1] We define a χ^2 statistic:

$$\chi^2\equiv-2\ln\mathcal{L}+\text{const}=\sum_{ij}\left[\ln\sigma_{ij}^2+\frac{S_{ij}^2}{\sigma_{ij}^2}+2\frac{(b_i\delta t_j-D_{ij})S_{ij}}{\sigma_{ij}^2}+\frac{(D_{ij}-b_i\delta t_j)^2}{\sigma_{ij}^2}\right]. \tag{12}$$

The model is nonlinear, and the variance σ_{ij}^2 depends on the model parameters, so we have defined this quantity including the $\ln\sigma_{ij}$ term; the usual χ^2 distribution does not exactly apply.

Because of the two terms contributing to the variance, the form of χ^2 depends on which term dominates. For $S_{ij}\gg b_i\delta t_j$, $\sigma_{ij}^2\approx S_{ij}$, and the S_{ij}^2/σ_{ij}^2 term dominates, so $\chi^2\sim\Sigma S_{ij}$. When the neutrino signal is small, the background contribution dominates, and $\chi^2\simeq$const. These regimes are shown in Fig. 2, where we plot χ^2 as a function of $m_\nu\tau$ for several values of B_γ.

Immediately, we see the character of the limits on the parameters. For $m_\nu\tau\lesssim10^7$ keV sec, $\chi^2\propto B_\gamma$; in this regime only the data from the first time bin after the supernova contributes. For $m_\nu\tau/B_\gamma\lesssim10^{13}$ keV sec, $\chi^2\propto B_\gamma/m_\nu\tau$; now, the full data set provides information. Finally, $\chi^2=$const for $m_\nu\tau/B_\gamma\gtrsim10^{13}$ keV sec; in this regime the background dominates over the theoretical signal. Note that this latter area of parameter space provides the "maximum likelihood" (or χ^2 minimum); there is no neutrino signal and we calculate only limits on parameters. In fact, there is a slight deficit of counts with respect to the background calculated from the time before the supernova; otherwise, we might expect to see a weak maximum likelihood somewhere in the large-$m_\nu\tau$ regime. In Fig. 3, we show a single contour of χ^2 in the $m_\nu\tau$-B_γ plane.

[1] We include a small contribution reflecting the uncertainty in that rate; this latter effect is somewhat more difficult to include if a Poisson distribution is explicitly used but is in any case negligible.

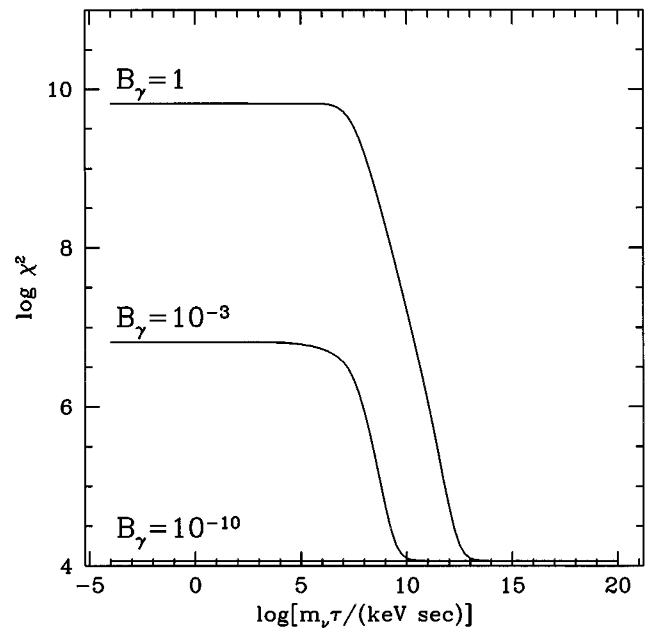

FIG. 2. The value of the χ^2 statistic, defined in the text, as a function of the parameter $m_\nu\tau$, for values of the branching ratio B_γ as marked, for the decay process $\nu\rightarrow\nu'\gamma$. The logarithm is to base 10.

To connect with other analyses, we shall assume that our χ^2-like statistic has a χ^2 distribution. We have 550 time bins and four detectors, so there are 2200 degrees of freedom. For this distribution, a 1σ fluctuation corresponds to $\Delta\chi^2=2230$, a 3σ (or 99%) fluctuation to $\Delta\chi^2=2357$; we choose the latter as our limit; from the shape of the likelihood function it is clear that any comparable $\Delta\chi^2$ will give similar bounds. We also note that a signal in the small-$m_\nu\tau$ regime may not be detectable with this algorithm; the absence of a local minimum in that region, however, implies that this should not be a significant worry. The allowed region is shown in Fig. 3. It corresponds approximately to

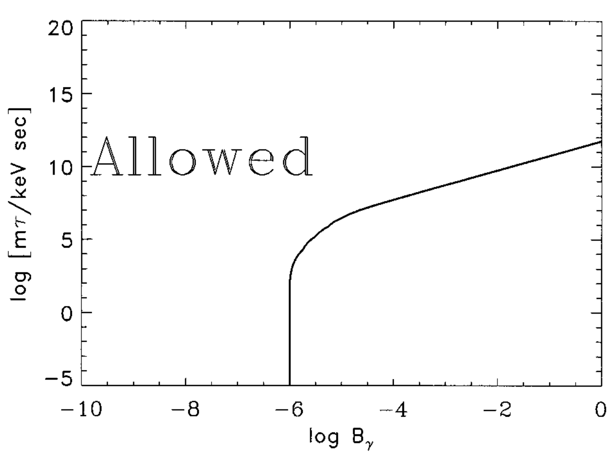

FIG. 3. Allowed region of the $m_\nu\tau$-B_γ plane, for the 2-body decay process $\nu\rightarrow\nu'\gamma$, corresponding to $\Delta\chi^2\leqslant 2357$ (see text). Here and below, contours continue to infinity as long as the appropriate assumptions, discussed in the text, still hold. The logarithm is to base 10.

$$B_\gamma<3\times10^{-7},\quad m_\nu\tau\lesssim10^6\ \text{keV sec},$$

$$B_\gamma<2\times10^{-13}\frac{m_\nu\tau}{\text{keVsec}},\quad m_\nu\tau\gtrsim10^6\ \text{keV sec},\tag{13}$$

for neutrinos with a temperature of 8 MeV, appropriate for μ and τ neutrinos. (The limits scale roughly as T^{-2}.) This is less restrictive than the limits of Oberauer *et al.* [6], due to the fact that the PVO GBD could only detect γ rays with energies below 2 MeV, compared to 25 MeV for the SMM Satellite. However, we believe this analysis to be more rigorous and the PVO data to be of higher quality.

IV. OTHER DECAY MODES—BREMSSTRAHLUNG

So far we have considered only photons produced from the simplest radiative decay mode of a massive neutrino species: $\nu\rightarrow\nu'\gamma$; three- and four-body modes are also possible. For these modes, where the rest-frame photon energy is no longer given by the simple $\overline{k}=m_\nu/2$, we must allow for a distribution of decay products: $f(\overline{k},\overline{\mu})d\overline{k}\,d\overline{\mu}$ gives the fraction of photons produced with rest-frame energy $\overline{k}$ into angle $\overline{\mu}=\cos\overline{\theta}$. Then, the final spectrum is

$$\frac{dN}{dkdt}=\frac{1}{4\pi D^2}\frac{B_\gamma}{\tau}\int dEL_\#(E)\times\int d\overline{\mu}f\left[\frac{k}{\gamma(1+v\overline{\mu})},\overline{\mu}\right]e^{-\gamma(1+v\overline{\mu})t/\tau},\tag{14}$$

where we have still assumed that the decays occur near the supernova and $m_\nu<T$, and the energy integral is taken from k to ∞. (We discuss the case of large masses, $m_\nu\gtrsim T_\nu$ in Sec. V below.) In particular, we consider the bremsstrahlung process, $\nu\rightarrow\nu' e^+e^-\gamma$, where $\nu=\nu_\tau$, $\nu'=\nu_e$. Because this is no longer a two-body decay, the kinematics are considerably more complicated, and no exact calculation of the spectrum has been performed. Following Oberauer *et al.* [6], we make several simplifications: (1) We assume isotropy of the photons in the rest frame of the neutrinos; this is reasonable if the helicity states of the parent neutrinos are produced in equal numbers. This gives $f(\overline{k},\overline{\mu})=f(\overline{k})/2$ (which still implicitly depends on $\overline{\mu}$ through the Lorentz transformation to the laboratory frame). (2) Up to factors of order unity, we assume that the bremsstrahlung energy spectrum is given by [6]

$$\frac{d\Gamma_{\rm br}}{d\overline{k}}\equiv\frac{B_\gamma}{\tau}f(\overline{k})\simeq\frac{\alpha}{\pi}\frac{\Gamma_0}{\overline{k}}=\frac{\alpha}{\pi}\frac{1}{\overline{k}\tau_e},\tag{15}$$

where Γ_0 and τ_e refer to the process without a daughter photon: $\nu\rightarrow\nu' e^+e^-$, absorbing a branching ratio factor into $\tau_e=\tau/B_e$. Now, the γ-ray flux is

$$\frac{dN}{dkdt}=\frac{1}{4\pi D^2}\frac{1}{m_\nu\tau_e}\frac{\alpha}{\pi}\int dEL_\#(E)\frac{E}{k}\times\int_{-1}^{+1}\frac{d\overline{\mu}}{2}(1+v\overline{\mu})e^{-\gamma(1+v\overline{\mu})t/\tau};\tag{16}$$

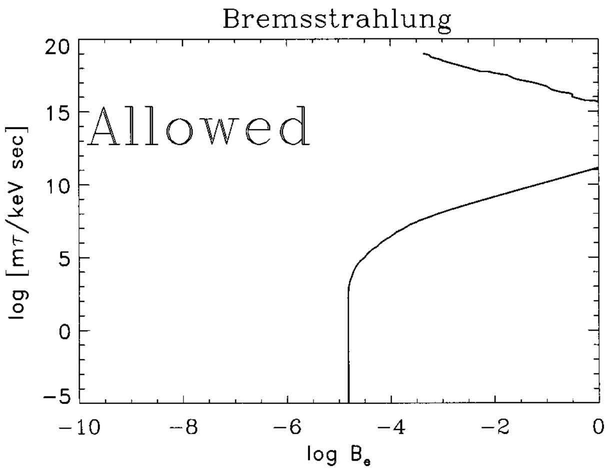

FIG. 4. Allowed region of the $m_\nu\tau$-B_e plane, for the bremsstrahlung decay process, as above. The logarithm is to base 10.

integrating over decay angle and the time bin gives a flux at the detector

$$\phi(k,t)=\frac{1}{4\pi D^2}\frac{\alpha}{\pi}\frac{1}{m_\nu\tau_e}\frac{1}{2k}\int dE L_\#(E)E\left(\frac{m_\nu\tau}{Et}\right)^2 \times[1-e^{-2Et/m_\nu\tau}(1+2Et/m_\nu\tau)]. \quad (17)$$

The integration over the time bin can be done analytically, but the remaining energy integration must be computed numerically. [For $\gamma t\ll\tau$, the angular average becomes unity and using a Maxwell-Boltzmann distribution for $L_\#(E)$ reproduces Eqs. (6) and (7) of Oberauer *et al.* [6].] With the usual assumption that a negligible fraction of the decays occurs inside of the progenitor and that detected γ rays have energies $k\ll T$ (i.e., $E_{\min}\ll T$), the flux is again a function only of $m_\nu\tau$. For $m_\nu\tau\ll T\delta t$, where δt is the timing resolution of the detector, the flux is proportional to $B_e/m_\nu\tau=1/m_\nu\tau_e$; for $m_\nu\tau\gg Tt$, it is proportional to $B_e m_\nu\tau$. This latter behavior results in an excluded region with $B_e\propto 1/m_\nu\tau$. The photon energy dependence is $1/k$ (from the spectrum) times a slowly-varying function of k/T from the integration over neutrino energies. Again, we show the allowed region of parameter space in Fig. 4.

Using the same definition of a 99% confidence level gives limits of approximately

$$B_e<2\times10^{-5},\quad m_\nu\tau\lesssim10^6\ \text{keV sec},$$

$$B_e<7\times10^{-12}\frac{m_\nu\tau}{\text{keVsec}},\quad m_\nu\tau\gtrsim10^6\ \text{keV sec},$$

$$B_e<5\times10^{15}\left(\frac{m_\nu\tau}{\text{keV sec}}\right)^{-1}. \quad (18)$$

Because the bremsstrahlung spectrum peaks at a lower energy and the time baseline of the PVO data is much longer, this limit is comparable to other SN 1987A limits for this channel [5,6], and, we believe, more reliable, due to the more careful calculation of the flux and the higher-quality data set.

V. VERY MASSIVE NEUTRINOS

All of these expressions are considerably more complicated in the case $m_\nu\gtrsim T$. We will still assume a zero-chemical-potential FD distribution, this time applying to massive particles:

$$L_\#(E,m_\nu)=\frac{120}{7\pi^4}\frac{E_T}{T_\nu{}^4}j(m_\nu/T_\nu)\frac{E\sqrt{E^2+m_\nu{}^2}}{1+e^{E/T_\nu}}, \quad (19)$$

with a "suppression factor" $j(x)$, along with the requirement that $E>m_\nu$. The factor $j(x)$ is just the usual Boltzmann suppression [$j(x)\propto x^{3/2}e^{-x}$, for $x\gg1$]; we use an approximation good for $m/T\lesssim$ few, $j(x)\simeq\exp(-0.15x^2)$. (Sigl and Turner [2] have calculated the effect of the changing neutrinosphere temperature and radius on this naive expectation; the effect is small for $m_\nu\lesssim40$ MeV, and $\tau\gtrsim10^{-2}$ sec.) For low-mass neutrinos we assumed $k>m_\nu$; now, we can only integrate over neutrino energies greater than max(m_ν,k). In this expression, $E_T\simeq10^{53}$ erg remains the total energy for the low-mass case; the total energy released is $E_T j(m_\nu/T_\nu)$, which for large masses is less than 10^{53} erg since $j<1$.

In addition to the mass threshold effects, we must now take into account the loss of photons produced inside the envelope of the supernova, $R_{\text{env}}=100c$ sec$\approx3\times10^{12}$ cm. Thus, we require that $vt_d>R_{\text{env}}$, or

$$E>E_{\text{env}}=m_\nu\sqrt{\frac{1+(R_{\text{env}}m_\nu)}{2kt}}. \quad (20)$$

Note that $E_{\text{env}}>m_\nu$, so this supersedes the requirement that $E>m_\nu$, but the requirement that $E>k$ remains. Thus, we must integrate over neutrino energies from $E_{\min}=\max(k,E_{\text{env}})$. This integration, the equivalent of $h(k/T)$ above, cannot be done in closed form, but again it can be approximated by a Gaussian [at least for the isotropic case $f(E,k)=1/p$]:

$$\frac{dN}{dkdt}=\frac{B_\gamma}{4\pi D^2}j(m_\nu/T_\nu)\frac{120}{7\pi^4}\frac{E_T}{T_\nu{}^2}\frac{2k}{m_\nu\tau}e^{-2kt/m_\nu\tau}g(E_{\min}/T), \quad (21)$$

with

$$g(x)=\int_x^\infty dy\frac{x}{1+e^x}\simeq\frac{\pi^2}{12}e^{-0.2x^2}; \quad (22)$$

the factor of 0.2 in the exponent approximates the shape of the integral for $x\lesssim$ few. This differs from the massless case by a total suppression factor

$$\frac{j(m_\nu/T_\nu)g(E_{\min}/T)}{h(k/T)}\simeq\exp\left[-0.15\left(\frac{m_\nu}{T_\nu}\right)^2-0.2\left(\frac{E_{\min}}{T_\nu}\right)^2\right]. \quad (23)$$

Since $E_{\min}\geq m_\nu$, this is always less than $\exp[-0.35(m_\nu/T_\nu)^2]$ for interesting masses $m_\nu\gtrsim T_\nu$; unfortunately, the time now appears in the expression for $E_{\min}$, so the dt integral is no longer trivial. First, then, let us consider the suppression factor if we ignore the effect of decays inside the supernova envelope, integrating from $E_{\min}=\max(m_\nu,k)$.

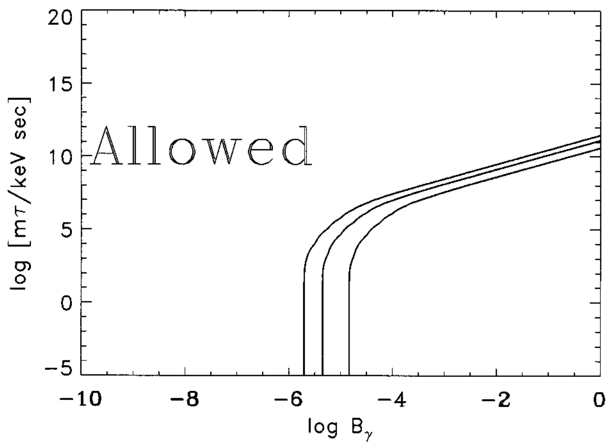

FIG. 5. Allowed region of the $m_\nu\tau$-B_γ plane, for neutrinos of mass 20 MeV, 30 MeV, and 40 MeV (from left to right), for the 2-body decay process, corresponding to $\Delta\chi^2 \leq 2357$ (see text). The logarithm is to base 10.

Then the time integral can be done as in the low-mass case, and we can simply write down the time-independent suppression factor $s = j(m_\nu/T) g(E_{\min}/T) \simeq j(m_\nu/T) g(m_\nu/T)$.

These mass effects enable us to break the degeneracy between m_ν and τ, at the price of a more complicated analysis of a three-dimensional parameter space. To simplify matters, we will base the results for massive neutrinos directly on the limits from the low-mass case. That is, we will calculate the limits as before, and then apply the suppression factor at the end. We can do this because the suppression factor comes into the expression for the flux in exactly the same way as the branching ratio B_γ, so we translate limits on B_γ in the massless case to limits on $B_\gamma \times s$, where s is the k-independent part of the suppression factor. In addition, we do the calculation for an isotropic decay, and assuming $k < m_\nu$. For two-body decay, this results in the limit

$$sB_\gamma < 3 \times 10^{-7}, \quad m_\nu\tau \lesssim 10^6 \text{ keV sec},$$

$$sB_\gamma < 6 \times 10^{-14} \frac{m_\nu\tau}{\text{keV sec}}, \quad m_\nu\tau \gtrsim 10^6 \text{ keV sec}. \tag{24}$$

If we allow the effect of decays inside the progenitor envelope, the calculation is somewhat more complicated, and the integration over each time bin can no longer be done in closed form. We must now recompute everything at each pair of m_ν and τ. We show the results of such a calculation for several values of the neutrino mass in Fig. 5; the limits are not too different from those with the simpler time-independent suppression.

For the bremsstrahlung process, the suppression of a high-mass neutrino flux is simpler to calculate because the required integral is simply $\int dE E L_\#(E)$, the total energy in the massive neutrino species [cf., Eq. (16); we have again assumed $k \ll T$, so the initial integration over μ simplifies]. Again, we integrate from the same $E_{\min} = \max(k, E_{\text{env}})$; the suppression is given by the Boltzmann factor $j(E_{\min}/T)$. For much of parameter space, this is simply the expected $j(m_\nu/T)$. As before, the time dependence of E_{env} does not change the limits significantly. The allowed parameter space for the bremsstrahlung process with a neutrino mass of 30 MeV is

$$\frac{m_\nu\tau_e}{j(m_\nu/T)} > 1.5 \times 10^{12} \text{ keV sec}$$

or

$$j(m_\nu/T) \frac{B_e}{m_\nu\tau} < 7 \times 10^{-13} \text{ keV}^{-1} \text{ sec}^{-1}. \tag{25}$$

VI. NEARLY DEGENERATE NEUTRINOS

Thus far, we have assumed that the daughter neutrino in the $\nu \to \nu' \gamma$ channel is much less massive than the parent neutrino. If, however, the mass of the daughter is appreciable, the energy of the photon will be decreased by a factor $\delta m^2/m^2 \equiv (m_1^2 - m_2^2)/m_1^2$. For $1 > \delta m^2/m^2 \gtrsim 1/100$, this improves our limits, shifting the bulk of the photons down from energies too high to detect into one or more of the energy channels of the PVO detector. To make the matter more precise, we see that in the case of nearly degenerate neutrinos, we make the change

$$f(E,k)dk \to f[E,(m^2/\delta m^2)k](m^2/\delta m^2)dk, \tag{26}$$

where we now are constrained to have photon energies $k < (\delta m^2/m^2)E$. This, in turn, results in changing $h(k/T) \to (m^2/\delta m^2) h[(m^2/\delta m^2)k/T]$. As expected, the flux is enhanced by as much as $(m^2/\delta m^2)$, along with another factor accounting for the shift of the spectrum. For $1 \lesssim \delta m^2/m^2 \lesssim 1/10$ essentially all of the photons fall in the PVO detector bands. For massless daughter neutrinos, of order 1/10 of the photons can be detected; therefore, we might expect limits as much as an order of magnitude stronger. In fact, for this case, the lower energy window of the PVO detector (down to 0.1 MeV) compared with the SMM window (sensitive only above 4.1 MeV) is actually an advantage. Unfortunately, for smaller values of $\delta m^2/m^2 \lesssim 1/100$, many or most of the daughter photons have energies below the PVO detector window, and remain unobserved (although they would perhaps be detectable as x rays). For $\delta m^2/m^2 \ll 1$, the flux is enhanced by approximately $[\delta m^2/m^2]^{-2} \exp[-(k/T_\nu) m^2/\delta m^2]$ over the nondegenerate case.

In Fig. 6, we show the limits on the neutrino parameters for $\delta m^2/m^2 = 10^{-3}$, as well as a more favorable (that is, detectable) possibility of $\delta m^2/m^2 = 0.1$; for values considerably below 10^{-3}, very few photons would be seen by even the lowest-energy detectors of PVO (of course, SMM, sensitive to even higher energies, would fare even worse).

VII. LONG LIFETIMES

For long lifetimes (such that the average decay time of the neutrino is comparable to or longer than the travel time to the detector), the above formalism becomes too cumbersome, because we must integrate over a complicated set of possible paths for the neutrino and daughter photon. In this case, we will make several simplifications. At first, we will only concern ourselves with the total γ-ray fluence from the decays,

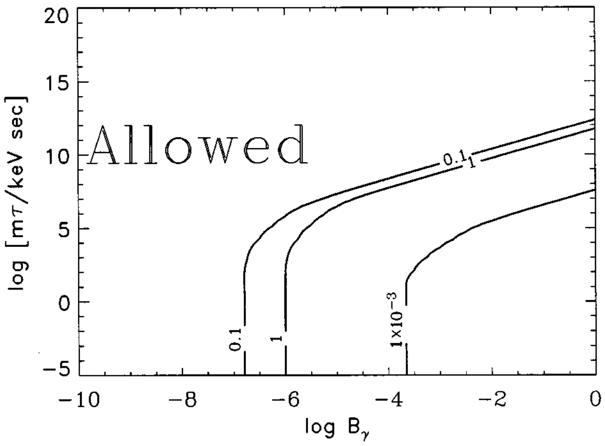

FIG. 6. Allowed region of the $m_\nu\tau$-B_γ plane, for nearly degenerate neutrinos, with $\delta m^2/m^2=10^{-3},0.1,1$ as labeled. The logarithm is to base 10.

integrated over time. Then, we will integrate over the decay time, $0\leqslant t_d\leqslant D$:

$$\frac{dN}{dk}=\frac{B_\gamma}{4\pi D^2}\int dEL_\#(E)f(E,k)[1-e^{-Dm_\nu/E\tau}] \quad (27)$$

where, as before, $f(E,k)$ gives the fraction of neutrinos with energy E decaying into photons with energy k. This expression is to be compared with those presented in Ref. [5]. The cost of the simplicity of this expression is the inability to determine the exact time of a photon's arrival. For neutrinos and photons traveling on a straight path ($\mu=1$, appropriate for relativistic particles), the arrival time after the supernova light pulse is $t=t_d(1-v)\simeq(t_d/2)m_\nu{}^2/E^2$. For long lifetimes, we will be concerned with neutrinos that decay late in their flight: $t_d\sim D$. Using this as a typical e-folding time, we have the ansatz that $(dN/dt)\propto\exp(-2tE^2/Dm_\nu{}^2)$. For $t\gtrsim(D/2)m_\nu{}^2/E^2$, this should express the character of the time dependence. Two effects are explicitly missing from this formula: the extra time delay from nonstraight paths (of the same order as the delay already considered) and the photon energy dependence of the time delay. Moreover, the time dependence will not have exactly this shape; for short times it does not contain the expected slow rise from zero flux, so it is probably safest to use this formula integrated over the entire duration of the experiment, and not rely on the detailed time evolution. We are left finally with

$$\left.\frac{dN}{dk}\right|_{\delta t}\simeq\frac{B_\gamma}{4\pi D^2}\int dEL_\#(E)f(E,k)(1-e^{-2\delta tE^2/Dm_\nu{}^2})\times(1-e^{-Dm_\nu/E\tau}). \quad (28)$$

Unfortunately, the integration over neutrino energy E is considerably more complicated than before, but we can approximate the two exponential decays for various regimes:

$$(1-e^{-2\delta tE^2/Dm_\nu{}^2})(1-e^{-Dm_\nu/E\tau})\approx\begin{cases}1, & m\lesssim E\sqrt{2\delta t/D}, \quad m_\nu/\tau\gtrsim E/D,\\ 2\delta tE^2/Dm_\nu{}^2, & m\gtrsim E\sqrt{2\delta t/D}, \quad m_\nu/\tau\gtrsim E/D,\\ Dm_\nu/E\tau, & m\lesssim E\sqrt{2\delta t/D}, \quad m_\nu/\tau\lesssim E/D,\\ 2\delta tE/m_\nu\tau, & m\gtrsim E\sqrt{2\delta t/D}, \quad m_\nu/\tau\lesssim E/D.\end{cases} \quad (29)$$

Numerically, these breaks occur at

$$m\simeq E\sqrt{2\delta t/D}\simeq 680\ \text{eV}\frac{E}{12\ \text{MeV}}\left(\frac{\delta t}{8500\ \text{sec}}\right)^{1/2},$$

$$m_\nu/\tau\simeq E/D\simeq 2\times10^{-6}\ \text{eV/sec}\frac{E}{12\ \text{MeV}}, \quad (30)$$

where $E=12$ MeV is a typical energy for a $T_\nu=8$ MeV blackbody. In terms of the lifetime τ, the latter limit occurs at $\tau\simeq Dm_\nu/E\approx5\times10^5\,\text{sec}(m_\nu/\text{eV})$—for masses $m_\nu\sim1$ MeV, this is roughly $\tau\sim D$. When this is proportional to δt, the flux is approximately constant; otherwise, the entire pulse is detected (and its shape is irrelevant). Putting all of this together, and doing the integration over E, gives

$$\left.\frac{dN}{dk}\right|_{\delta t}\approx\frac{B_\gamma}{4\pi D^2}\frac{120}{7\pi^4}\times\begin{cases}\dfrac{E_T}{T_\nu^2}h_0(k/T),\\ 2\dfrac{E_T}{m_\nu{}^2}\dfrac{\delta t}{D}h_2(k/T),\\ \dfrac{E_Tm_\nu}{T_\nu^3}\dfrac{D}{\tau}h_{-1}(k/T),\\ 2\dfrac{E_T}{T_\nu m_\nu}\dfrac{\delta t}{\tau}h_1(k/T),\end{cases} \quad (31)$$

where $h_n(y)=\int_y^\infty x^{n+1}/(1+e^x)$ is similar to $h(y)$ above, and the four cases correspond to those in Eq. (29). Here, we have assumed an isotropic distribution of decays in the rest frame. As before, these expressions hold for $m_\nu\lesssim T_\nu$ and must be modified with the appropriate suppression factor otherwise.

Now, we can just put these results through our statistical machinery and find limits on the parameters. We will write the flux as

$$\left.\frac{dN}{dk}\right|_{\delta t}\approx\frac{1}{4\pi D^2}\frac{120}{7\pi^4}\frac{E_T}{T_\nu^2}h(k/T)\times B_\gamma A(m_\nu,\tau) \quad (32)$$

where A is the appropriate dimensionless combination of m_ν and τ, along with D, T_ν and δt; the data give us limits on A in each (m_ν,τ) regime. This gives an approximate 99% confidence limit of $B_\gamma A\lesssim1\times10^{-6}$ for $T_\nu=8$ MeV or

$$B_\gamma \lesssim 1\times 10^{-6}\left(\frac{T_\nu}{8\ \mathrm{MeV}}\right), \qquad m_\nu \lesssim 0.4\ \mathrm{keV}, \quad m_\nu/\tau \gtrsim 1.2\times 10^{-9}\ \mathrm{keV/sec},$$

$$m_\nu \gtrsim 155\ \mathrm{keV} B_\gamma^{1/2}, \qquad m_\nu \gtrsim 0.4\ \mathrm{keV}, \quad m_\nu/\tau \gtrsim 1.2\times 10^{-9}\ \mathrm{keV/sec},$$

$$\frac{B_\gamma m_\nu}{\tau} \lesssim 1.4\times 10^{-15}\ \mathrm{keV\,sec}^{-1}\left(\frac{T_\nu}{8\ \mathrm{MeV}}\right)^3, \qquad m_\nu \lesssim 0.4\ \mathrm{keV}, \quad m_\nu/\tau \lesssim 1.2\times 10^{-9}\ \mathrm{keV/sec},$$

$$m_\nu \tau \gtrsim 1.4\times 10^{14}\ \mathrm{keV sec} B_\gamma\left(\frac{T_\nu}{8\ \mathrm{MeV}}\right)^{-1}, \qquad m_\nu \gtrsim 0.4\ \mathrm{keV}, \quad m_\nu/\tau \lesssim 1.2\times 10^{-9}\ \mathrm{keV/sec}. \tag{33}$$

Where the regimes overlap, these limits are comparable to those calculated with the more detailed models above; because we can only calculate limits on parameters, the details of the data and the analysis are unimportant [in fact, the limits of Eq. (33) are *stronger* than, for example, Eq. (13) above; the earlier, more detailed calculation is probably the more appropriate limit]. Again, for neutrinos with $m_\nu \gtrsim T_\nu$, these limits are modified with $B_\gamma \to s B_\gamma$.

For the bremsstrahlung channel, the flux is changed due to the different kinematics of the decay [i.e., the rest-frame spectrum of Eq. (15)]:

$$\left.\frac{dN}{dk}\right|_{\mathrm{brem}} = \frac{\alpha}{\pi}\frac{2T_\nu^2}{k m_\nu}\times\left.\frac{dN}{dk}\right|_{\text{two body}} \tag{34}$$

(in addition, the functions h_n should also be modified to h_{n+2}). This is a significant increase in flux at for $k m_\nu \lesssim T^2$. As before, we see that the bremsstrahlung spectrum at the detector is proportional to $1/k$. Now, the limits correspond to $B_\gamma A T/m \lesssim 3\times 10^{-5}$ or

$$m_\nu \gtrsim 2.7\times 10^8\ \mathrm{keV} B_\gamma\left(\frac{T_\nu}{8\ \mathrm{MeV}}\right), \qquad m_\nu \lesssim 0.4\ \mathrm{keV}, \quad m_\nu/\tau \gtrsim 1.2\times 10^{-9}\ \mathrm{keV/sec},$$

$$m_\nu \gtrsim 800\ \mathrm{keV} B_\gamma^{1/3}\left(\frac{T_\nu}{8\ \mathrm{MeV}}\right), \qquad m_\nu \gtrsim 0.4\ \mathrm{keV}, \quad m_\nu/\tau \gtrsim 1.2\times 10^{-9}\ \mathrm{keV/sec},$$

$$\tau/B_\gamma = \tau_e \gtrsim 1.9\times 10^{16}\ \mathrm{sec}, \qquad m_\nu \lesssim 0.4\ \mathrm{keV}, \quad m_\nu/\tau \lesssim 1.2\times 10^{-9}\ \mathrm{keV/sec},$$

$$m_\nu \gtrsim 1.4\times 10^8\ \mathrm{keV}\left(\frac{\tau/B_\gamma}{\mathrm{sec}}\right)^{-1/2}\left(\frac{T_\nu}{8\ \mathrm{MeV}}\right), \qquad m_\nu \gtrsim 0.4\ \mathrm{keV}, \quad m_\nu/\tau \lesssim 1.2\times 10^{-9}\ \mathrm{keV/sec}. \tag{35}$$

VIII. DISCUSSION

SN 1987A not only confirmed astrophysicists' standard model of Type II (core collapse) supernovae, but also because of the high fluence of neutrinos (about $10^{11}\ \mathrm{cm}^{-2}$ per species) provided a laboratory for studying the properties of neutrinos. This large fluence and the space-borne γ-ray detectors operating on SMM and PVO have allowed stringent limits to be placed on the radiative decay of neutrinos.

Although there are only 232 sec of data, in a single time bin, the SMM detectors were sensitive up to energies of 25 MeV, and are better matched to decaying thermal neutrinos with a temperature of 4–8 MeV. Because the limits are determined by the region of parameter where the background becomes comparable to the signal (see the discussion in Sec. III), the long time base and greater resolution is actually of little use in improving the limits on the parameters. To show this, we have performed our analysis with the SMM data, as well; as expected, the results are comparable to those calculated in Refs. [5,6]. In the case where the γ rays are produced by bremsstrahlung or the neutrino mass states are nearly degenerate, the PVO limits are even much more stringent. Finally, the amount and quality of the PVO data adds additional confidence to the SMM-based limits.

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