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A LOOK AT SUPERNOVA 1987A*

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ABSTRACT

Supernova 1987A is reviewed with emphasis on the neutrino observations. It is shown that the results fit well with the expectations for neutrino temperatures ($T \sim 4\epsilon_0 4.5 MeV$) and total energy emitted ($2\epsilon_0 4 \times 10^{53} ergs$). It is argued that the detection tends to favor collapse models that yield emission for 10 second timescales with a

$1\epsilon_0 2$ second early accretion phase followed by Kelvin-Helmholtz cooling as opposed to prompt shocks with the immediate onset of cooling. It is also argued that the probable detection of one or more electron scattering event favors a superthermal tail at high energies. Neutrino mass limits and flavor limits are comparable to laboratory experiments. An estimate for future collapse rates in our galaxy of 1/7 year is made based on nucleosynthesis yields. The supernova also has eliminated many axion and majoron models.

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Contrary to the attempt by the popular press to imply mystery, SN 1987A is actually reasonably well understood. In particular, many stellar evolution calculations (c.f. Lamb, Iben and Howard[1], 1977; Truran and Brunish[2], 1975) had indicated that massive stars could die with their outer regions as blue rather than red stars either because they were never red or because they were once red but contracted to the blue either by mass loss or simple contraction of the convective zone. Thus the fact that the progenitor was Sanduleak-69 202, a blue supergiant of 15 to 20 $M_\odot$, is not surprising. Once it is known to be a compact blue star then the low luminosity, high velocity ejection naturally follows. (Blue means higher temperatures, thus smaller radius for a given stellar luminosity. Smaller radius at the time of shock breakout yields lower initial supernova luminosity. Similarly, the dense compact blue envelope yields higher velocities as the shock propagates down the density gradient.)

The traditional concerns of Type I vs Type II are only concerns of whether or not hydrogen was present in the outer envelope, and for traditional Type II's it has to be spread out in red giant form. The real physics question is whether the star was massive enough to undergo core collapse or whether the explosion was caused by a nuclear detonation (deflagration) of a C-O white dwarf. Collapse yields detectable neutrinos, detonation does not. Thus this was a core collapse, the light curve questions are more analogous to surface weather.

For over 20 years, it has been known that the gravitational collapse events, thought to be associated with Type II supernovae and neutron star or black hole formation, are copious producers of neutrinos. In fact, the major form of energy transport in these objects comes from neutrino interactions. It has long been predicted that the neutrino fluxes produced by these events would be high enough that if an event occurred within the galaxy, it could be detected. The following discussion of the neutrinos borrows heavily from Schramm (1987)[3].

It has been well established in the models of Arnett[4] and Weaver et al.[5] that massive stars with $M \gtrsim 8\mathrm{M}_\odot$ evolve to an onion-skin configuration with a dense central iron core of about the Chandrasekhar mass surrounded by burning layers of silicon, oxygen, neon, carbon, helium, and hydrogen. Collapse inevitably occurs when no further nuclear energy can be generated in the core.

Bethe and Brown[6] and Baron et al.[7] have argued that, provided the equation of state of matter above nuclear density is very soft, stars in the mass range $10 \leq M \lesssim 16\mathrm{M}_\odot$ may explode due to the prompt exit of the shock wave formed after the core bounces upon reaching supra-nuclear density. For stars with $16 \lesssim M \leq 80\mathrm{M}_\odot$, the shock wave stalls on its exit from the core and becomes an accretion shock. Wilson et al.[8] have shown that such stars will eventually (~ 1 second later) eject their envelope as a result of neutrino heating in the region above the neutrinosphere and below the shock. (The delayed ejection can also occur in the lower mass collapses if the initial bounce does not produce an explosion.) In fact, if collapse to a black hole is delayed by about a second after bounce, the neutrino spectra and mass ejection should not be affected by the later formation of the black hole. Obviously the above scenarios are sensitive to the stiffness of the core equation of state which is still poorly known at and above nuclear mass densities. As was first emphasized by Arnett and Schramm[9], the ejecta have a composition which fits well with the observed 'cosmic' abundances for the bulk of the heavy elements.

Regardless of the details of collapse, bounce, and explosion, it is clear that to form a neutron star the binding energy, $\epsilon_B \approx 2 \times 10^{53}$ ergs must be released. The total light and kinetic energy of a supernova outburst is about 10^{51} ergs. Thus, the difference must come out in some invisible form, either neutrinos or gravitational waves. It has been shown[10] that gravitational radiation can at most carry out 1% of the binding energy for reasonable collapses because neutrino radiation damps out the non-sphericity of the collapse (see Kazanas and Schramm[11,12]). Thus, the bulk ($\gtrsim 99\%$) of the binding energy comes off in the form of neutrinos.

It is also well established[13] that for densities greater than about 2×10^{11} g/cm^3, the core is no longer transparent to neutrinos. Thus, as Mazurek[14] first established, the inner core has its neutrinos degenerate and in equilibrium with the matter. For electron neutrinos, the 'neutrinosphere' has a temperature such that the average neutrino energy is around 10 MeV. This was established once it was realized that the collapsing iron core mass is $\sim 1.4\mathrm{M}_\odot$, due to the role of the Chandrasekhar mass in the pre-supernova evolution. Since the μ and τ neutrinos and their antiparticles only interact at these temperatures via neutral rather than charged current weak interaction, their neutrinosphere is deeper within the core. Therefore, their spectra are hotter than the electron neutrino. The electron antineutrino opacity will initially be dominated by charged current scattering off protons but as the protons disappear, it will shift to neutral current domination. Thus the effective temperature for $\bar{\nu}_e$'s changes from that for ν_e's to that for ν_μ and ν_τ's.

The average emitted neutrino energy is actually quite well determined (Schramm 1987, ref.3, also ref.15) for the peak of the neutrino distribution and is very insensitive to model parameters. The peak occurs at the highest temperature for which neutrinos can still free stream out of the star; that is, where the neutrino mean free path, $[n\langle\sigma\rangle]^{-1}$, is comparable to the size of the core, R. This can be expressed as

$$R \simeq 1/n\langle\sigma\rangle \tag{1}$$

where n is the number density $= \rho/m_N$, ρ being the mass density and m_N the nucleon mass. Collapsing stars are well described by adiabatic physics. Thus density and temperature are related as

$$\rho = \rho_0 (T/T_0)^3 \tag{2}$$

For a Fermi distribution the average energy $\langle E_\nu\rangle = 3.15T_\nu$ (using T in energy units). The effective neutrino cross section in stars[16] can be expressed as

$$\langle\sigma\rangle \approx \sigma_0\langle E_\nu^2\rangle/2 \approx 12\sigma_0/2T_0^2\left(\frac{T_\nu}{T_0}\right)^2$$

Inserting this into eq. 1 and solving for T

$$\frac{T_{\nu_e}}{T_0} = [\frac{2m_n}{R\rho_0 12\sigma_0 T_0}]^{1/5} \tag{3}$$

The neutrino temperature, T, varies only as the 1/5 power of the input. Thus, large uncertainties get minimized. [If R is put at its upper limit from the size of the core, then $R \propto 1/T$. The limiting relationship has T proportional to the 1/4 power, which is still quite insensitive.] Using reasonable values $\sigma_0 \simeq 1.7 \times 10^{-44}$ and $\rho_0 = 10^{10}$ gm/cm^3 at $T_0 = 1$ MeV, with the characteristic size of the region $R \lesssim 5 \times 10^6$ cm (see models[5,6]). Then

$$T_{\nu_e} \approx [340]^{1/5}\text{MeV} \simeq 3.2\text{MeV}$$

or

$$\langle E_\nu\rangle \approx 10\text{MeV}.$$

This is in good agreement with detailed numerical results. As mentioned above, for $\bar{\nu}_e$'s the average energy increases with time. The time averaged value is about 15MeV.

It should also be noted that since the interaction cross sections in the star are proportional to the square of the neutrino energy, the lower energy neutrinos can escape from deeper in the star. In addition, as time goes on the core evolves, so some higher energy neutrinos are able to get out from deeper inside. Thus, the energy distribution of the emitted neutrinos is not a pure thermal distribution at the temperature of the neutrinosphere.

While the general scenario for collapse events is well established, the detailed mechanism for the ejection of the outer envelope in a supernova as the core collapses to form a dense remnant continues to be hotly debated. Therefore, most theorists working on collapse prior to SN 1987A have focused on these details in an attempt to solve the mass-ejection problem. As a result, most of the pre-1987 papers in the literature are concerned with the role played by neutrinos internal to the stellar core, rather than the nature of the fluxes which might be observed by a neutrino detector on earth. In particular, while it has been known since the early 1970's[17,18] that the average energy of the emitted neutrinos was about 10 MeV, with neutrino luminosities of a few 10^{52} ergs/sec, the detailed nature of the emitted spectra was only recently explored in detail by Mayle, Wilson, and Schramm[19,20]. Their calculation emphasized the high-energy neutrinos which are easier to detect. The diffusion approximation used in most collapse calculations does not treat the high-energy tail of the spectrum accurately. A large temperature gradient exists in the neutrinospheric region. For the high-energy neutrinos, the matter's temperature at one optical depth is relatively low compared to the temperature at one optical depth of the mean-energy neutrino. Thus, an appreciable fraction of the high-energy neutrinos originate in the higher temperature region and travel several mean-free paths before exiting the star. Therefore, for neutrinos whose energy is far above the mean energy, the multi-group, flux-limited-diffusion approximation is suspect. To confirm this, Mayle et al. constructed a computer code that integrates the Boltzmann equation more directly. Calculations by other authors using diffusion are fine as far as energy transport in the core is concerned but breakdown on questions of detailed neutrino spectra.

In addition to the basic energetic arguments, there is the neutronization argument (see ref. 3, and references therein). The collapsing core has $\sim 10^{57}$ protons that are converted to neutrons via

$$p + e^- \rightarrow n + \nu_e$$

to form a neutron star. (This process is also called deleptonization by some authors.) Each ν_e, so emitted from the core, carries away on the average 10 MeV, thus around 1.3×10^{52} ergs are emitted by neutronization ν_e's. this is $\lesssim 10\%$ of the binding energy. The remainder of the neutrinos come from pair processes such as

$$e^+ + e^- \rightarrow \nu_i \bar{\nu}_i$$

where $i = e$, μ, or τ, with ν_μ and ν_τ production occurring via neutral currents, and ν_e via both charged and neutral currents. (See review by Freedman et al.[21].)

Since some fraction ($\lesssim 50\%$) of neutronization occurs in the initial collapse, whereas the pair ν's come from the 'thermally' radiating core, the timescale for an initial neutronization ν_e burst will be much less ($\lesssim 10^{-2}$ sec) than the diffusion time ($\sim$ seconds) that governs the emission of the bulk of the flux. Some so-called 'advection/convection' models increase the initial ν_e burst by convecting high-T, degenerate core material out. These models have higher-energy ν_e's with larger fluxes, and suppress the $\bar{\nu}_e$ fluxes.

Even in the "detailed" explosion models, more than half of the thermal neutrino emission comes out in the first one or two seconds with prompt models having even a greater fraction emitted in less time. The remainder comes out over the next few tens of seconds as the hot, newborn, neutron star cools down via Kelvin-Helmholz neutrino cooling to become a standard 'cold' neutron star. Burrows and Lattimer[22] carried out detailed cooling calculations prior to SN 1987A. Most other authors cut off their calculations after the bulk of the neutrino emission occurred and mass ejection was established. Detailed models for the bulk of the neutrino emission (c.f. Mayle et al.[19]) seem to find that the pair processes yield an approximate equipartition of energy in the different species. The ν_μ and ν_τ's have a higher energy per ν, thus their flux is down to preserve this equipartition.

Despite the explosive mechanism, for stars in the mass range $10 \lesssim M \lesssim 16\mathrm{M}_\odot$ the most distinctive structure in the neutrino signal is the initial neutronization burst. However, in the delayed explosions seen by Wilson et al.[8], for stars with $M \geq 16\mathrm{M}_\odot$, besides the burst, the neutrino luminosity shows an oscillatory behavior superimposed on an exponentially decaying signal. The oscillations in luminosity are related to oscillations in the mass accretion rate onto the proto-neutron star. The physical nature of the instability that is responsible for the oscillations in luminosity and mass-accretion rate is described in Wilson et al.[8], and in more detail in Mayle[23]. After the envelope is ejected, the luminosity will smoothly decrease as the remaining binding energy is emitted during the Kelvin-Helmholz cooling. Models without the accretion phase go directly from the neutronization burst to Kelvin-Helmholz cooling. Those models thus have the $\bar{\nu}_e$ emission fall off with a single characteristic cooling time. However, models with an accretion phase have a high average emission rate for a second or so after the neutronization burst before the mass ejection and onset at the cooling phase with its dropping emission.

It is important to remember that the average neutrino luminosity, mean neutrino energy, and total emitted energy depend only on the initial iron-core mass and are otherwise independent of the explosive mechanism. Because the opacity is less for the ν_μ and ν_τ's, they are emitted from deeper in the core where temperature is higher. Thus, they have a higher average energy. The calculations of Mayle et al.[19] find $E_{\nu_\mu} \simeq E_{\nu_\tau} \approx 2E_{\nu_e}$. The easier-to-observe $\bar{\nu}_e$'s start out with energy comparable to ν_e's and gradually shift over to the $\nu_\mu - \nu_\tau$ energy as their emission continues from progressively deeper in the core.

Each spectrum for neutrino species is reasonably well fit by a Fermi–Dirac (F–D) distribution with temperature T. However, in the detailed spectral analyzer that Mayle et al.[19] carried out, it was found that the higher-energy neutrino fluxes were indeed higher than the single-temperature, F–D fit to the peak.

We'll wait until we analyze the individual detectors before discussing sensitivities to thresholds, etc. However, just by using simple, model-independent arguments, one obtains a crude $\bar{\nu}_e$ counting rate for an H_2O detector

$$n = \frac{(1-f_n)\epsilon_B}{2N_\nu \langle E_\nu \rangle} \frac{\langle \sigma \rangle}{4\pi r^2} \frac{2}{18} \frac{M_D}{m_p} \tag{4}$$

where f_n is the fraction radiated in the neutronization burst, $\langle E_\nu \rangle$ is the average neutrino energy, $\langle \sigma \rangle$ is the average cross section above threshold. [It should be noted that the cross section goes as peEe not E_ν^2, (see discussion Appendix to reference 3). However, this effect can be treated as an additional detector sensitivity factor.] r is the distance to the LMC $\approx$ 50 Kpc, M_D is the mass of the detector, m_p is the proton mass, and N_ν is the number of neutrino flavors. (For the Mt. Blanc liquid-scintilator detector, one should multiply by 1.39 for the average number of free protons in $H_{2+2n}C_n$.) Using F–D statistics yields

$$\langle \sigma \rangle \approx \frac{\int_{E_c}^{\infty} \bar{\sigma} \frac{E^4 dE}{1+e^{E/T}}}{\int_0^{\infty} \frac{E^2 dE}{1+e^{E/T}}} \tag{5}$$

where E_c is the low-energy cut-off and $\bar{\sigma} \equiv \sigma/E_\nu^2$. Later, we'll discuss E_c and trigger efficiencies, however, for now let us do the crude estimate that everyone did before real data existed. Namely, let $E_c \to 0$, then

$$\langle\sigma\rangle \approx 7.5 \times 10^{-44} 12 T_{\bar{\nu}_e}^2 \mathrm{cm}^2 = 12\bar{\sigma} T_{\bar{\nu}_e}^2$$

Prior to SN 1987A, estimates were made for distances within our galaxy. With the LMC these had to be scaled by r^2. For completeness, let us plug in the standard numbers, $\epsilon_B = 2 \times 10^{53}$ ergs, $N_\nu = 3$, $f_n = 0.1$, and $T_{\nu_e} \sim 4$ MeV. Thus,

$$n = 5.2 \left(\frac{T}{4\mathrm{MeV}}\right) \left(\frac{\epsilon_B}{2 \times 10^{53}\mathrm{ergs}}\right) \left(\frac{1-f_n}{0.9}\right) \frac{1}{(N_\nu/3)} \left(\frac{M_D}{\mathrm{ktons}}\right) \left(\frac{50\mathrm{kpc}}{r}\right)^2 \quad (6)$$

For the 2.14 kiloton Kamioka detector, this yields 11 counts. Similarly, for the Mt. Blanc detector with 0.09 kilotons, times 1.39 extra, free protons in the scintillator, a simple prediction is $\sim$ 0.6 counts. IMB is a little more difficult because its threshold is not below the peak $\bar{\nu}_e$ counting rate. In addition, it is totally dominated by the high T tail where a constant T may not be an ideal approximation. However, we can crudely estimate that $\sim$ 50% of the $\bar{\nu}_e$ counting rate is above the approximate IMB low E cut-off of 20 MeV. Thus, with 5 kilotons, IMB should roughly get 13 effective counts. If we are more careful regarding efficiencies' thresholds and integrals over F–D distributions we reduce this prediction to 6. However, even the crude estimates show about what one naively expected from supernova theory independent of detailed models.

To estimate the expected number of electron scattering events one must do a bit more if threshold effects are to be included. Electron scattering yields a very flat energy distribution. When such a flat energy distribution is combined with a finite temperature F-D distribution for the initial neutrinos, one finds an expected energy distribution for the scattered electrons which is quite peaked at low energies. If pure constant temperature F-D distributions are assumed for the neutrinos, the total number of scattering is expected to be $\lesssim$ 0.5 for $10\bar{\nu}_e$ capture events. If the high energy tails are supressed by absorption as Imshennik and Nadyoshen (c.f. 24 and references therein), then the expected scattering rate is even lower. However, if the high energy super-thermal tails of Mayle et al. are included, one finds that for every 10 $\bar{\nu}_e$ absorptions, one expects about 0.7 to 1 ν_e scattering and about 0.7 $\nu_x e$ scattering, where ν_x is either ν_μ, $\bar{\nu}_\mu$, ν_τ, $\bar{\nu}_\tau$, or $\bar{\nu}_e$. We can understand why the scattering rate is

$\sim 1/15$ even though the cross section ratio at 10 MeV is ~ 80 by remembering that there are five electrons for each free proton in an H_2O target. In addition, at a given energy from our cross section table

$$(\sigma_{\nu_\tau^e} + \sigma_{\nu_\mu^e} + \sigma_{\bar{\nu}_\tau^e} + \sigma_{\bar{\nu}_\mu^e} + \sigma_{\bar{\nu}_e^e})/\sigma_{\nu_e^e} \simeq 1.$$

Thus, if fluxes are equal, the rate is doubled. Actually, average energy of other species is about twice that of ν_e, but fluxes are reduced accordingly to roughly maintain equipartition of energy per neutrino species, thus keeping scattering constant. The difference in expected number of scatterings is an important probe of the high energy tail.

For the 615-ton C_2Cl_4 Homestake there are 2.2×10^{30} ^{37}Cl atoms. As seen from the Appendix to reference 3, the cross section is not a simple integer power of E_ν, however, it seems to fall roughly between E^3 and E^4 relationship for $E_\nu \lesssim 30 MeV$. For temperatures above 5 MeV, the peak contribution to the thermal average would be coming from energies above 30 MeV where the cross section no longer rises as rapidly and the expected counting rate no longer continues to rise with temperature. In the standard case, one expects about a half of a count above the background. However, for advection models, one might expect several ^{37}Cl events. Similar to the solar case, ^{37}Cl is once again a potentially sensitive thermometer.

All the predictions described above assume a simple, spherical symmetric collapse. If large amounts of rotation or magnetic fields were present (with energies comparable to the binding energy) then the standard model would be altered with different time scales and different core masses and binding energies, since such conditions would alter the initial core mass as well as the dynamics. We will see that the Kamioka/IMB neutrino burst fits the standard assumptions well so that the collapse which created that burst did not have significant rotation or magnetic fields.

Before SN 1987A, it was also obvious that a supernova, if detected by its neutrinos, would constrain neutrino properties. In particular, if the neutrinos got here, we'd have a lifetime limit. If the time pulse wasn't too spread out, that would mean a mass limit on those neutrino types that were clearly identified. Also, from the number of $\bar{\nu}_e$ counts, one could constrain N_ν since if N_ν was large, the fraction of thermally produced $\bar{\nu}_e$'s would go down. In addition, neutrino mixing could be constrained by detecting different types and comparing; with

Mikheyev–Smirnov[16] matter mixing, as parameterized to solve the solar neutrino problem, $\nu_e \rightarrow \nu_\mu$ (or ν_τ), and ν_μ (or ν_τ) $\rightarrow \nu_e$, but nothing happens in the antineutrino sector. Such mixing would eliminate seeing the initial ν_e burst, but give higher energies to the later, thermal ν_e since they'd be mixed ν_μ's (see Walker and Schramm[25]). Of course, non-solar Mikheyev–Smirnov can be used if antineutrino mixing is seen. All of these effects will be examined with the data from SN 1987A.

NEUTRINO OBSERVATIONS

Table 1 summarizes the neutrino observations, noting two reported neutrino bursts. Before discussing the plausibility of the first event, it is important to note that *all* neutrino detectors clearly had a detection on February 23rd near 7h 35m U.T. Thus, unquestionably *extra solar system neutrino astronomy has been born!* Let us now examine the burst Mt. Blanc reported on February 23rd, -2:52 with five events which was unsubstantiated by the other three detectors. While lack of concordance is easy to understand for IMB and Baksan, due to their higher thresholds, the lack of a strong concordant signal, significant above background, is difficult with regard to Kamioka. The Kamioka detector is 2140 tons, compared to 90 tons for Mt. Blanc. (Mt. Blanc was designed to detect $\bar{\nu}_e$'s from collapses in our galaxy, not the LMC.) Thus, many people have dismissed this first event as an unfortunate statistical accident. *A posteriori* statistics are difficult. While the chance of background exactly duplicating this event configuration eight hours before the visual outburst is low, perhaps the more relevant question is: What is the chance of background producing any plausible signal within two days prior to the visual detection? If any plausible signal is defined as three or more events (only three events were clearly above background) in less than or equal to 30 seconds, a chance occurrence becomes quite reasonable and many have assumed this explanation. However, one should be cautious in following popular opinion too rapidly. Detections near threshold can be tricky, and statistics of small numbers are notoriously suspect. In fact, while both thresholds are indeed low, Mt. Blanc is lower. In particular, Mt. Blanc sees positrons down to $\sim$ 5MeV whereas Kamioka does not see positrons below 7MeV (their 50% efficiency point is actually 5 MeV). Furthermore, Mt. Blanc sees total energy including etc. e^+e^- annihilation thus is capable of detecting incoming neutrinos down to 5.3 MeV whereas Kamioka must add 1.3 MeV to get their neutrino energies, yielding their lower bound on detectable neutrinos of 8.3 MeV, a full 3 MeV above Mt. Blanc.

Kamioka did report that they had a background count in the 10-minute interval centered at the Mt. Blanc event which is consistent with their background. Although preliminarily De Rujula[26] had thought that if the IMB burst is used to accurately set the Kamioka U.T. clock (which was only calibrated to $\sim \pm 1$ minute absolute), and the Kamioka background is scanned at precisely the U.T. of the first Mt. Blanc burst, then the high count Kamioka sees at the 10-minute interval happens to fall within eight seconds of the Mt. Blanc event. However, recently the Kamioka team (Koshiba, priv. Comm.) has carried out this exercise in detail and found that their event is about 20 seconds away from the Mt. Blanc events. Thus Kamioka appears to have no concordant events with Mt. Blanc at the early time. Table 2 shows the implied temperature and neutrino luminosity implied by the Mt. Blanc burst and the one or two Kamioka counts at that time. These were estimated by deconvoluting F–D distributions with thresholds and efficiencies. Notice that the burst reported at Mt. Blanc is not well fit by the standard collapse assumptions but instead requires lower-than-expected temperatures and extraordinarily high total energies.

Let us suspend our theoretical prejudice and ask if such a high-luminosity, low-T event did occur, could Kamioka not have seen it? In fact, as first noted by De Rujula[26] a minimal Kamioka detection cannot be totally excluded because the implied Mt. Blanc burst temperature is so low, and the thresholds are different. Even zero events is possible if the temperature of the neutrino distribution were low enough. To get less than a few counts at Kamioka requires neutrino temperatures under 1 MeV. Lower temperatures yield higher flux in order to get 5 events at Mt. Blanc. To avoid a Kamioka conflict would require $T \leq 1$ MeV and $E_{TOTAL} \gtrsim 10^{55} ergs$! Some have also cited "$3\sigma$" gravitational wave detector noise in Italy and Maryland in coincidence with the Mt. Blanc burst as significant. However, these are room temperature detectors with lots of noise and would imply $> 2000 M_\odot$ emitted in gravitational waves at LMC. The Mt. Blanc burst would necessitate an initial collapse event that is quite different from standard models. Models with large magnetic fields and/or rotation, such as Symbalisty et al.[27] have low temperatures, but it is hard to imagine an event which radiates a *minimum* of several neutron star *rest masses* in neutrinos, or has a very non-thermal distribution. The non-standard event must then be followed by a subsequent collapse five hours later to a black hole or a dense, strange-matter star looking very much like a normal collapse, as we shall see. An alternative is that this event was not in the LMC but was *much* closer, thus reducing the energy requirements but requiring a remarkable timing coincidence. Given all these problems, we quote Eddington: "Observations should not be believed until confirmed by theory". Unlike Kamioka and IMB, it should be remembered that the Mt. Blanc detector was actually constructed to look for collapse neutrinos; unfortunately it was optimized for collapses within 10 kpc.

Let us now turn our attention to the well established Kamioka/IMB burst. (For a detailed discussion, the fact that Mt. Blanc and Baksan also have signals is irrelevant other than to show that detectors $\sim 1/20$ the mass can have $\sim 1/10$ the counts, due to statistics of small numbers plus possible background subtraction uncertainties.) Figure 1 is a plot showing the energy and timing of the Kamioka and IMB events. (Kamioka's event no. 6 is ignored as being below their criteria for a definitive event.) Note that almost all the counts concentrate in the first few seconds, as one expects in collapse models. A reasonable tail, as predicted by theory[19,27], yields low but finite rates after ten seconds. Such rates following the bulk early emission from an accretion phase have little difficulty in producing apparent gaps in counts due to the problems of small number statistics (c.f., Bahcall[28] et.al. or Mayle and Wilson[29]). Note also that the IMB late counts nicely fill in the 6 second gap in the Kamioka data. Prompt explosion models with cooling starting almost immediately have some trouble fitting both the initial two second high counting rate and the long time tail accretion phase (Mayle et al.) followed by cooling have no problem. Unfortunately small number statistics make it difficult to make categorical statements.

To examine consistency let us use the number of counts and mean energies measured in the experiments to determine the implied temperature and energy emitted in $\bar{\nu}_e$'s. Such estimates require detailed consideration of efficiency and threshold effects.

To convert a mean neutrino energy to an effective temperature requires assuming that the emitted ν spectrum was well described by Fermi–Dirac statistics. Mayle et al. argue that this is a reasonable assumption, however, as mentioned before, they did find that their models had a higher tail at high energies than a simple, single-temperature model would yield. Thus, one might expect the IMB temperature to be slightly higher than the Kamioka temperature due to its weighting on the high-energy events. If the $\bar{\nu}$'s fit F–D statistics, then the mean energy $\langle E_\nu \rangle$ as recorded by a detector with cross section proportional to E_ν^2 and cut-off energy E_0, with efficiency of detection $f(E)$, is given below, where E and T are measured in the same units, and $E_0 = E_c(e^+) + Q$

$$\langle E_\nu \rangle = \frac{\int_{E_0}^{\infty} \frac{f(E)E^5 dE}{1+e^{E/T}}}{\int_{E_0}^{\infty} \frac{f(E)E^4 dE}{1+e^{E/T}}} \simeq \frac{E_0 + 5T + 20\frac{T^2}{E_0} + \frac{60T^3}{E_0^2} + \frac{120T^4}{E_0^3} \frac{120T^5}{E_0^4}}{1 + \frac{4T}{E_0} + \frac{12T^2}{E_0^2} + \frac{24T^3}{E_0^3} + \frac{24T^4}{E_0^4}} \tag{7}$$

which goes to the well known F–D integral values for $E_0 = 0$. Thus, we have a polynominal equation for T:

$$T^5 + \left(E_0 - \frac{\langle E_\nu \rangle}{5}\right) T^4 + \left(\frac{E_0^2}{2} - \frac{\langle E_\nu \rangle E_0}{5}\right) T^3 + \left(\frac{E_0^3}{6} - \frac{\langle E_\nu \rangle E_0^2}{10}\right) T^2 + \left(\frac{E_0^4}{24} - \frac{\langle E_\nu \rangle E_0^3}{30}\right) T + \frac{E_0^5 - \langle E_\nu \rangle E_0^4}{120} = 0 \tag{8}$$

This latter equation can be trivially solved for the effective temperature, $T_{\bar{\nu}_e}(\langle E_\nu \rangle,\ E_0)$: from this equation it is obvious that the effective T is a very sensitive function of E_0. We will use efficiency weighted values for n so as to avoid the treatment of the efficiency function in the integrals. Evaluating $\langle \sigma \rangle$ counts and energies from equation 5 yields

$$\langle \sigma \rangle \approx \frac{7.5 \times 10^{-44}}{2} T_{\bar{\nu}_e}^2 \left(\frac{E_0^4}{2T^4} + \frac{2E_0^3}{T^3} + 6\frac{E_0^2}{T^2} + 12\frac{E_0}{T} + 12\right) e^{-E_0/T} \mathrm{cm}^2$$

again, a function that is sensitive to E_0. Equation 6 for n can be inverted to solve for $\epsilon_{\bar{\nu}_e}$ where the total energy, ϵ_T (which can be compared to neutron star binding energy, ϵ_B) is related to $\epsilon_{\bar{\nu}_e}$ by

$$\epsilon_T \approx \frac{2N_\nu \epsilon_{\bar{\nu}_e}}{(1 - fn)}$$

. The numbers in Figure 2 are calculated assuming $N_\nu = 3$ and $f_n = 0.1$, with Kamioka having $M_D = 2.14$ kilotons, and IMB having $M_D = 5$ kilotons. Figure 2 shows the energy radiated versus $T_{\bar{\nu}_e}$. The boundaries of the region come from one σ errors in counts as well as the range of reasonable assumptions one might make about cut-off energies and stated experimental errors in energy.

While one might expect (from Mayle et al.) IMB to measure a slightly higher T, it is interesting that there is nevertheless a region of overlap where both data sets yield the same T_{ν_e} and $\epsilon_{\bar{\nu}_e}$. It is particularly satisfying that this region of overlap is exactly where one might have expected a standard gravitational collapse event to plot, namely, $\epsilon_T \sim 2 \times 10^{53}$ ergs, $T \sim 4.5$ MeV. Similar conclusions were reached by Sato and Suzuki[30] and Bahcall et al.[28] using a different treatment than has been applied here. Once T and ϵ_T are determined one can use the luminosity–temperature relationship to solve for the radius, R, of the neutrinosphere and obtain, in our case, a few tens of kilometers in reasonable agreement with the standard models, whether of not the first two or the last three events from Kamioka are included. However, the high E_c Kamioka data set with minimal weighting effects does seem to yield those parameters which are closer to overlap with IMB and closer to expected supernova parameters. It is worth noting that the above analysis is very crude, Kolb et al.[31] have pointed out that simple converting of E_e to $E_\nu - Q$, as was done here, is inaccurate although it does not effect these conclusions. Also note that the boundaries used in Figure 2 do not have a quantitative statistical meaning since systematic as well as statistical uncertainties were mixed in obtaining them. Nonetheless, the results are suggestive and more detailed analyses seem to yield similar conclusions[28,29,30].

The IMB angular distribution initially was thought to be biased due to a failed power supply however Monte Carlos by the IMB team find their data to be solid but with so few counts it is hard to make strong statements, although it does show more counts towards LMC than away from it. The angular distribution for Kamioka is shown in Figure 3. It appears to show an isotropic distribution with a possible slight excess in the direction of LMC. From the isotropic rate background and the angular resolution, the number of excess directed events (note, Kamioka only explicitly claims two probable scatterings, but considering resolution, etc., we feel that our estimate is reasonable) is $\sim 3 \pm 1.8$. Since $\bar{\nu}_e + p$ would yield an isotropic distribution, the number of directed electron scattering events should be relatively small, as might be expected by the ratio of cross sections. As mentioned before Mayle et al.[19] expect ~ 1.5 for $12 M_\odot$ or 2 for their 15 $M_\odot$ model in reasonable agreement with the observations. One also expects that $\sim 50\%$ of these scattering events are higher energy ν_μ, ν_τ, $\bar{\nu}_\mu$, $\bar{\nu}_\tau$, or $\bar{\nu}_e$ events. This also fits well since the highest energy Kamioka events have $\cos\theta > 0.7$. It is also intriguing that the first two events had $\cos\theta$ closest to unity. Remember that the initial 0.01 sec neutrino burst is expected to be ν_e's with no $\bar{\nu}_e$'s. While two such scatterings might be excessive considering the cross section suppression (unless the ν_e flux is slightly enhanced by advection convection) statistics of two versus one are not worth arguing about and are not useful in confirming or denying one theory instead of another. It is interesting to note that models with no high energy tail would predict less than 1/2 a scattering event. Since the data seems to require 1 or more with 3 as a best fit, it is reasonable to argue that the data do lean towards models with high energy tails over models with pure constant T distributions and certainly models with absorption supressed tails run into difficulty.

While discussing ν_e scattering, its worth noting that the ^{37}Cl experiment of Davis was operating at the time of The Supernova, and counting began shortly after the light was observed. This experiment is only sensitive to ν_e's. After 45 days of counting, Davis saw one count, completely consistent with his normal counting rate[32]. As mentioned before, for a standard collapse one expects from the LMC event ~ 0.5 events in the Homestake Chlorine detector. However, if one interprets the Kamioka data as implying a large excess[33] of ν_e's, then one might have expected several ^{37}Cl counts. The lack of observed Cl counts argues that the ν_e flux is not in disagreement with standard predictions of $\sim 2 \times 10^{52}$ ergs of neutronization ν_e's, plus 3×10^{52} ergs of thermal ν_e's, all at $E_\nu \sim 10$ MeV ($T_{\bar{\nu}_e} \sim 3.5\,\mathrm{MeV}$). This constrains models[34,35] with 'advection' producing excessively large high-energy ν_e fluxes and reducing the $\bar{\nu}_e$ fluxes. As mentioned earlier, such models can predict at most about 5 ^{37}Cl counts. While extreme models with $T_{\nu_e} \gtrsim 5\,\mathrm{MeV}$ and $f_{\nu_e} \sim 1$ may be in difficulty, intermediate models with $T_{\nu_e} \lesssim 4\,\mathrm{MeV}$ and/or $f_{\nu_e} \lesssim 0.5$ are still allowed.

Another constraint on ν_e's comes from interactions with ^{16}O which would be backward peaked at high energy. No data shows any evidence for this.

Before leaving the neutrino data, it cannot be emphasized too strongly that statistics of small numbers are *dangerous* and one must be extraordinarily careful *not to overly interpret* all the bumps, wiggles, and time delays. Impatience in waiting for a collapse event in our galaxy with ~ 100 times the counting rate is unfortunately stimulating such detailed interpretations of the only data we have. Conclusions drawn in this way must be appropriately normalized.

CONSTRAINTS ON NEUTRINO PHYSICS

Independent of detailed collapse models, we can use the detection of neutrinos from SN 1987a in the Kamioka and IMB detectors to constrain neutrino properties.

Neutrino Lifetime

Obviously, if $\bar{\nu}_e$'s made it over 50 Kpc, they must have a lifetime τ such that

$$\gamma\tau \gtrsim 1.6 \times 10^5 \mathrm{yr}$$

where γ is the relativistic factor ($\gamma = E_\nu/m_\nu$). Of course, to have decay requires $m_\nu > 0$. Since γ for ν's from the sun is $\sim 1/10$, γ's from supernovae (assuming $m_{\nu_e} = m_{\bar{\nu}_e}$) this means that neutrino decay is not a solution to the solar neutrino problem unless one combines decay with special mixing assumptions[36].

Decay

An additional lifetime constraint comes for any neutrino decaying to photons. In particular if the decay occurred in the star it would effect the dynamics (Falk and Schramm 1977) as long as $\gamma \gtrsim 10^{-3}$ sec. Even if it occurred in flight it would yield γ-rays associated with the neutrino burst at unacceptable levels. Thus we know that ν_e's ν_μ's and ν_γ's are able to go for 1.6×10^5 years without decaying to γ's.

Neutrino Mass

Since the neutrino bursts were relatively narrow in timespread, despite the energies being spread out over a range of about a factor of two, it is obvious that there cannot be too significant of a neutrino rest mass. While the relationship between mass, timespread and energy is derived in freshman physics the world over, the key here is to decide which counts to use to get the time and energy spread, and to estimate what the intrinsic spread was in the neutrino burst in the absence of finite masses. It is these assumptions that have yielded more neutrino mass preprints than neutrino events observed. (Thus, we will not bother to reference them.)

Before discussing what we can say in a model-independent manner, it is important to emphasize that all we get model-independently is an upper limit on the mass, since it is certainly possible that the timespread is just due to intrinsic emission time, and not any mass effects. Thus, all papers claiming finite masses rather than upper limits are intrinsically model-dependent. In addition, since most, if not all, of the counts are $\bar{\nu}_e$'s, it is only reasonable to measure neutrino-mass limits for $m_{\bar{\nu}_e} = m_{\nu_e}$, not for any other neutrino species unless assumptions about mixing are made. (Of course anything else, like a fine-tuned photino, that interacts in H_2O with a rate similar to $\bar{\nu}_e$, and is produced in supernovae, would also be limited.)

Let us now plug some values into the standard relation for the mass implied by two particles of energy, E_1 and E_2, emitted at the same time, but arriving 50 Kpc away with a separation Δt.

$$m = 20\text{eV}\left(\frac{E_1}{10\text{MeV}}\right)\left[\frac{(\Delta t/10\text{sec})}{(r/50\text{Kpc})}\frac{(E_2/E_1)^2}{(E_2/E_1)^2 - 1}\right]^{1/2}$$

Model-independently, the simplest thing to do is to assume that the entire 13 sec spread of Kamioka was due to this effect. (IMB, with its higher energies, isn't able to constrain things as well.) Schramm[3] and Kolb[31] argue that with these assumptions alone it is really difficult to get limits much better than $M_{\bar{\nu}_e} \lesssim 30eV$. Once we admit that the supernova limit is comparable to the Zurich experimental limit[37] of $m_{\nu_e} < 20$ eV, the whole game becomes irrelevant, except for the curiousity that by having the supernova take place in LMC, the values come out very close to terrestrial laboratory measurements.

Alternative games of assuming two or more neutrino types of different mass run into the problem of low cross section for detection of all but $\bar{\nu}_e$. In addition, if the three late Kamioka events were a different neutrino with $m \sim 20$ eV, compared to the earlier burst with $m_{\bar{\nu}_e} \ll 20$ eV, one also has trouble understanding why these late events don't show any strong directional character, since they would then be electron-scattering events for either a $\nu_\mu + \bar{\nu}_\mu$ or $\nu_\tau + \bar{\nu}_\tau$. While it would be wonderful to have $m_{\nu_\tau} \approx 20$ eV, to give us the hot dark matter of the universe, this supernova cannot be used to prove it (or disprove it).

If specific models are assumed, slightly tighter limits can be obtained. For example, Abbott, De Rujula and Walker[38] using a diffusing neutrinosphere model obtain a 90% confindence limit of $m_{\bar{\nu}_e} \lesssim 18$ eV for the Kamioka events and Bahcall and Spergle[39] find $m_{\nu_e} \lesssim 16eV$ if all 19 events are used and the relative timing of Kamioka and IMB is optimized. However Mayle and Wilson[29] show that their models fit the data equally well with any $m_{\bar{\nu}_e} \lesssim 30$ eV!

Number of Neutrino Flavors, Axions and Majorons

A limit to the number of neutrino flavors (with $m_\nu \lesssim 10$ MeV), N_ν, can be derived[40,3,41] from observation of the supernova-produced $\bar{\nu}_e$'s. The argument is based on the fact that in an equipartition of emitted neutrino luminosities among all flavors, the more flavors, the smaller the yield per flavor. Since $\bar{\nu}_e$ is only one flavor, this means that a detection of $\bar{\nu}_e$'s tells you immediately that the dilution by flavor could not have reduced the luminosity of $\bar{\nu}_e$'s below detectability. We can do this in a couple of ways; for example, from our simple relation for the predicted number of $\bar{\nu}_e$ counts in an H_2O detector, compared with the number observed, N_{obs}, we can calculate N_ν.

$$N_\nu \leq 3 \left[\frac{5.2}{n_{obs}} \left(\frac{T}{4\text{MeV}}\right) \left(\frac{\epsilon_B}{2\times 10^{53}\text{ergs}}\right) \left(\frac{1-fn}{0.9}\right) \left(\frac{M_D}{\text{kton}}\right) \left(\frac{50kpc}{R_{LMC}}\right)^2 \right]$$

Using n_{obs} weighted by the detector-efficiency yields for Kamioka 16.5 ± 5 events (14.3 ± 4.3 if two events are electron scatterings). Putting in the deviations in the cross section from E_ν^2 only strengthens the limits.

$$N_\nu \leq (2 \pm 0.6) \left[\left(\frac{T_{\bar{\nu}_e}}{4\text{MeV}}\right) \left(\frac{\epsilon_B}{2\times 10^{53}\text{ergs}}\right) \left(\frac{1-fn}{0.9}\right) \left(\frac{50kpc}{r}\right)^2 \right]$$

(If two events are assumed to be electron scattering, the 2 goes to 2.3.) From the concordant temperature prejudice, we can estimate that T is good to better than 25%. Similarly, ϵ_B for $1.4\text{M}_\odot$ neutron-star models doesn't go over 3×10^{53}, independent of equation of state (or for an extreme limit with $1.6\text{M}_\odot$, we will also use 4×10^{53} ergs.) Obviously, $1 - f_n$ can't exceed unity. Allowing 10% uncertainty in r and putting in our extreme values yields

$$N_\nu < 6.3(9.1)$$

We hesitate to use this method with IMB data because of the need to be more careful with thresholds in $\langle\sigma\rangle$. An alternative technique is to use our explicit results for $\epsilon_{\bar{\nu}_e}$, as implied by the experimental detections. Since $\epsilon_{\bar{\nu}_e}$ was derived using detailed integrations of cross sections with cut-off energies, we don't have the cross section averaging uncertainty, implicit in the previous technique. Noting that with equipartition of energies

$$\epsilon_{\bar{\nu}_e} = \frac{(1-f_n)}{2N_\nu}\epsilon_B$$

we can solve for N_ν

$$N_\nu = \frac{(1-f_n)\epsilon_B}{2\epsilon_{\bar{\nu}_e}}$$

Fitting to the center of the IMB–Kamioka consistent range, we find $\epsilon_{\nu_e} = 3.5 \times 10^{52}$ ergs, for $\epsilon_B = 2 \times 10^{53}$ ergs, and $1 - f_n = 0.9$. This yields

$$N_\nu = 2.9$$

If we take the extreme low value for ϵ_{ν_e}, $1 - f_n$ of 1, and again allow ϵ_B to be 3×10^{53} ergs (4×10^{53}), we find the limit,

$$N_\nu < 5.5(7.3),$$

quite compatible with our more simply derived limit, ignoring thresholds and adjusting Kamioka data. This number is not as restrictive as cosmological bounds[42,43] but is comparable to current accelerator limits[44].

This argument can be used to limit any other sort of particle that might be emitted by the supernova and dilute the $\bar{\nu}_e$ energy share. Using the fact that axions can escape from the higher T central core even though neutrinos cannot, we (Mayle, et al.[45]) can further restrict axion coupling, $f_a \gtrsim 10^12 GeV$ exceeding current red giant limits[46] marginally and eliminating axions altogether since cosmological density arguments constrain $f_a \lesssim 10^{12}$ GeV. Similarly, Fuller et.al. have shown that this supernova tightly constrains majorons with Frieman[47] arguing that it may eliminate them altogether.

Neutrino Mixing

If neutrino mixing occurs between emission and detection, it can obviously alter things. If the mixing is simple vacuum oscillations and the mixing length is short compared to 50 Kpc, then the chief effect will be an increase in the average ν_e, and to a lesser extent $\bar{\nu}_e$, energy, due to the oscillations with the higher energy ν_μ's and ν_τ's. Since we only reliably detect $\bar{\nu}_e$'s, this energy enhancement would be difficult to resolve. While some supernova models may need such enhancements to understand the IMB counts, others such as Mayle et al. do not; thus, no definite statements on mixing can occur. (The possibility of the electron scattering events having high energy is also still in the noise.)

Let us now address the matter mixing such as Mikheyev and Smirnov, and Wolfenstein[16] (MSW) have proposed. Walker and Schramm[25] have applied this to stellar collapse scenarios. If this is indeed the solution to the solar neutrino problem, then only $\nu_e \leftrightarrow \nu_\mu(\nu_\tau)$ mixing is possible, not $\bar{\nu}_e \rightarrow \bar{\nu}_\mu(\bar{\nu}_\tau)$. Thus, the solar neutrino solution would not enhance $\bar{\nu}_e$ fluxes. It would deplete the initial neutronization burst. Since ν_μ cross sections are down by $\sim 1/6$, the possibility of seeing a neutronization scattering is significantly reduced. Thus, if the possible scatterings are real, standard adiabatic MSW is not the solution to the solar neutrino problem. However, proving that the first two events in an eleven event distribution are really scattering rather than isotropic background is fraught with statistical difficulties.

If we drop the solar neutrino solution and go to general MSW mixing, then we can mix $\bar{\nu}_\mu(\bar{\nu}_\tau)$ into $\bar{\nu}_e$, which might enhance the energy slightly, but would otherwise do little. No effect would occur for the electron scattering ν_e's. As in the case of vacuum oscillations, no definitive statement can be made.

COLLAPSE RATES

Over the last 1000 years there have been only 5 visual supernovae in the Milky Way Galaxy, implying at first glance a rate of 1/200 years. However, if we look at galaxies like our own, that is standard evolved spiral Sb and Sc galaxies, we find[3] in other galaxies of 1/15 to 1/20 years. Obviously our galaxy's low rate is probably the result of most of our galaxy being obscured from view by dust in the disk. In fact the 5 historical supernovae were all in our sector of the galaxy implying a <u>minimal</u> enhancement of a factor of 5 to 1/50 yr to include the entire disk volume. Now that we can detect collapses by neutrinos alone, we don't need to worry about the obscuration of our disk, so the rates in other galaxies where we sample their entire disk might be more relevant. However with neutrino detectors we only see Type II supernovae thus the rates quoted may be on the high side since these include all types ("neutrinoless" Type I's account for $\sim 1/3$ to $1/2$ of the supernovae by such direct counting of supernovae in these galaxies). Such direct counting of supernovae is fraught with uncertainties. For example SN 1987A would probably not have been included since it was so underluminous. If the fraction of blue star collapsing is only minimally related to metallicity then SN 1987A types could enhance the supernova for the high metallicity disk populations. It may even be that metallicity enhances the blue progenitor fraction as high mass loss rates might move more stars from red to blue prior to collapse. Of course, if the blue progenitors only occur in metal poor populations, SN 1987A would not alter the statistics for the Milky Way. Similarly, other underluminous collapses, such as Cassioppe A would not be detected in extragalactic surveys. Tammann[48] discusses many other difficulties.

An alternative approach is to do statistics on stellar types. The rate of formation of all stars $\gtrsim 8M_\odot$ is $\sim$1/8 yr using a Salpeter mass function and a constant star formation rate. All such stars presumably undergo collapse. Of course the Salpeter mass function is probably most uncertain for these more massive stars, and the assumption of a constant rate can be argued. Similar numbers can be obtained from pulsar formation rates with even larger uncertainties.

We do not know that from the 2% heavy element content of our galaxy and the assumption that $\gtrsim 1M_\odot$ of heavies is ejected per collapse that the $10^{11}M_\odot$ disk requires $\lesssim 2 \times 10^9$ ejections over the 15×10^9 yr history of the galaxy. Thus our average Type II rate is $\lesssim 1/7$ yr. Since our current rate of explosion is $\lesssim$ the average, this is certainly a good limit. Since the current best galactic evolution models seem to have roughly constant nucleosynthesis rates[49] this limit is also not a bad estimate and is in good agreement with the Salpeter rate estimate. The fact that it is higher than the rates observed in similar galaxies probably argues that many collapses are indeed underluminous and were missed in the surveys.

SUMMARY

This supernova in the LMC has proven to be one of the most exciting astrophysical events of the century. It has already taught us much about supernova physics and more should be forthcoming as heavy element spectra and the remnant come into view. We now know that blue as well as red stars collapse, that SN luminocities for blue progenitors are indeed lower than for red ones and velocities are higher.

The neutrinos from SN 1987A have proven that our understanding of the basic energetics of gravitational collapse was quite reasonable once we included neutral current effects. Given that we now know what a neutrino burst looks like, we should have confidence that if a collapse occurs anywhere in our galaxy, regardless of the visibility of the SN, we should observe it. From SN rates in other galaxies like ours, we expect a rate of a collapse every 20 years or so and the neutrino flux will be up by $1/r^2$.

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FIGURE CAPTIONS

Figure 1: The ν counting rates for IMB/Kamioka.

Figure 2: Emitted energy, $\epsilon_{\bar{\nu}_e}$ in $\bar{\nu}_e$ and total emitted energy, ϵ_T (assuming $N_\nu = 3$) versus temperature for Kamioka and IMB data, allowing for statistical errors as well as systematic shifts due to possible electron scattering events and variations in threshold and efficiency assumptions. Note overlap region is a good fit to the standard model.

Figure 3: Angular distribution of Kamioka data. If level of isotropic events is chosen from directions away from LMC, then there appears to be $\sim 3 \pm 1.8$ excess counts in the direction of the LMC, presumably due to electron scattering. The standard model predicts ~ 1.5 to 2; pure F–D constant T predict < 0.5. Mayle et al.

FIGURE 1

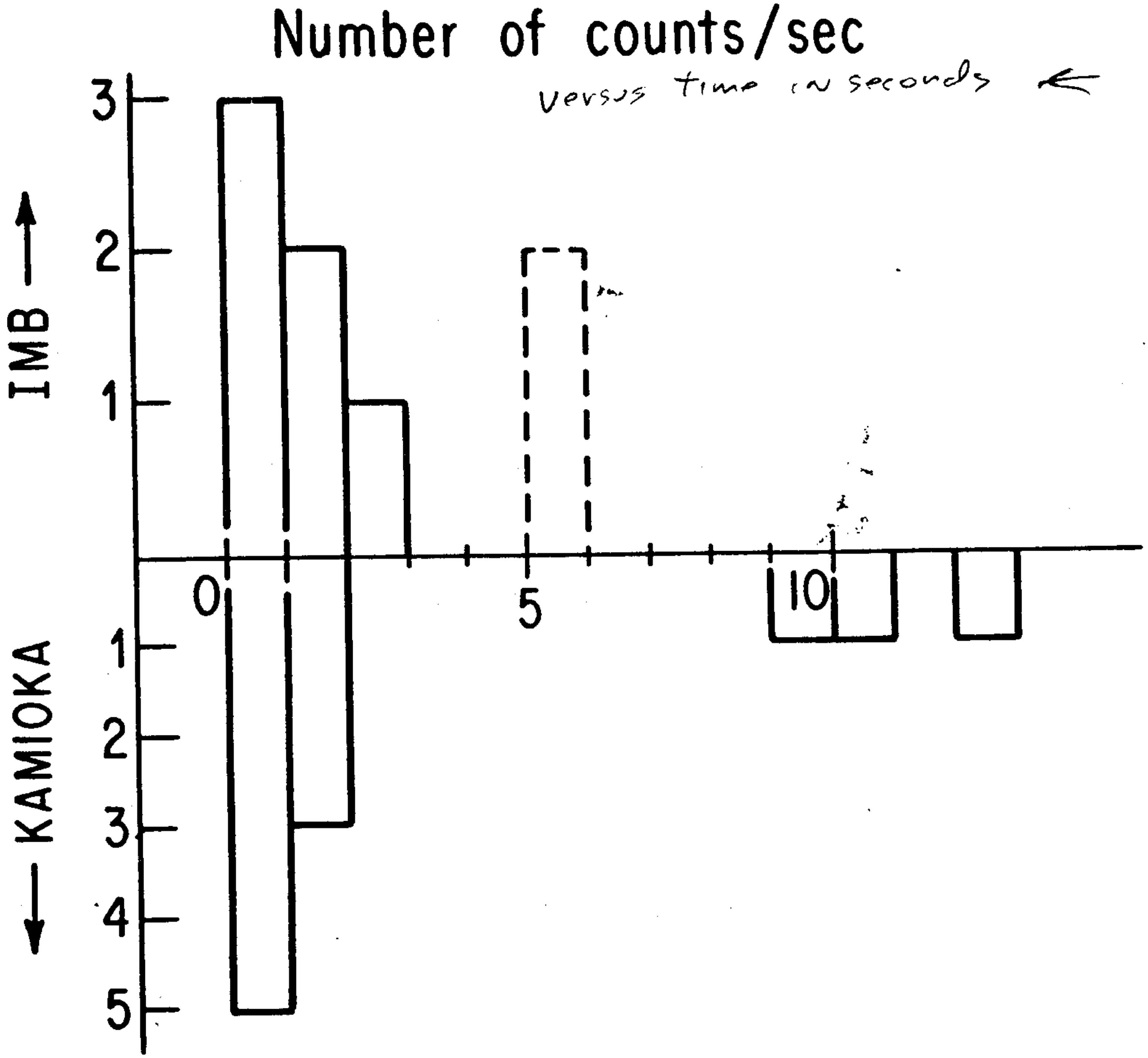
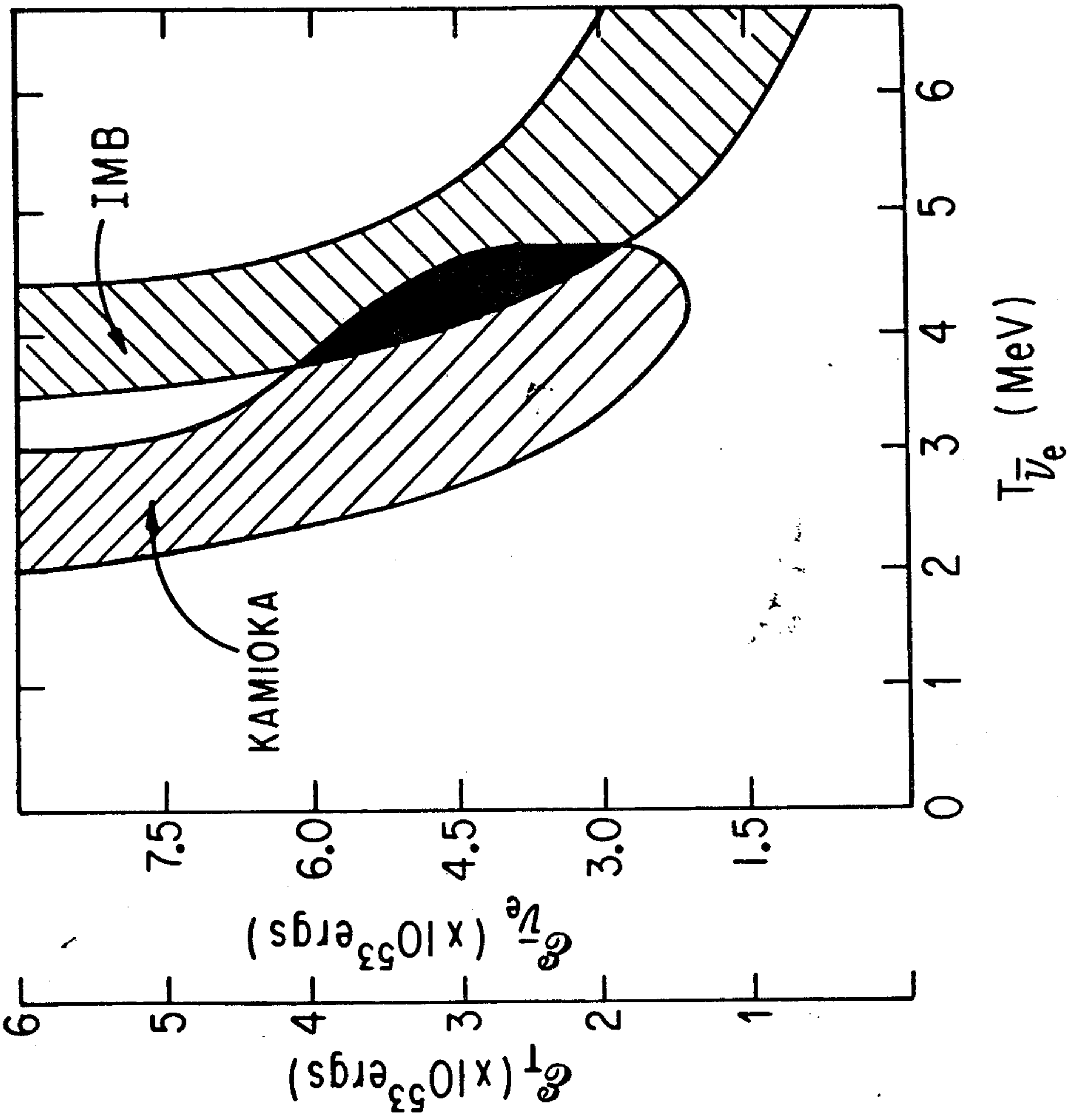

Number of counts/sec
versus time in seconds
IMB
KAMIOKA
3
2
1
0
5
10
1
2
3
4
5

FIGURE 2



IMB
KAMIOKA
$T_{\bar{\nu}_e}$ (MeV)
0 1 2 3 4 5 6
$\mathcal{E}_{\bar{\nu}_e}$ (x 10^{53} ergs)
1.5 3.0 4.5 6.0 7.5
$\mathcal{E}_T$ (x 10^{53} ergs)
1 2 3 4 5 6

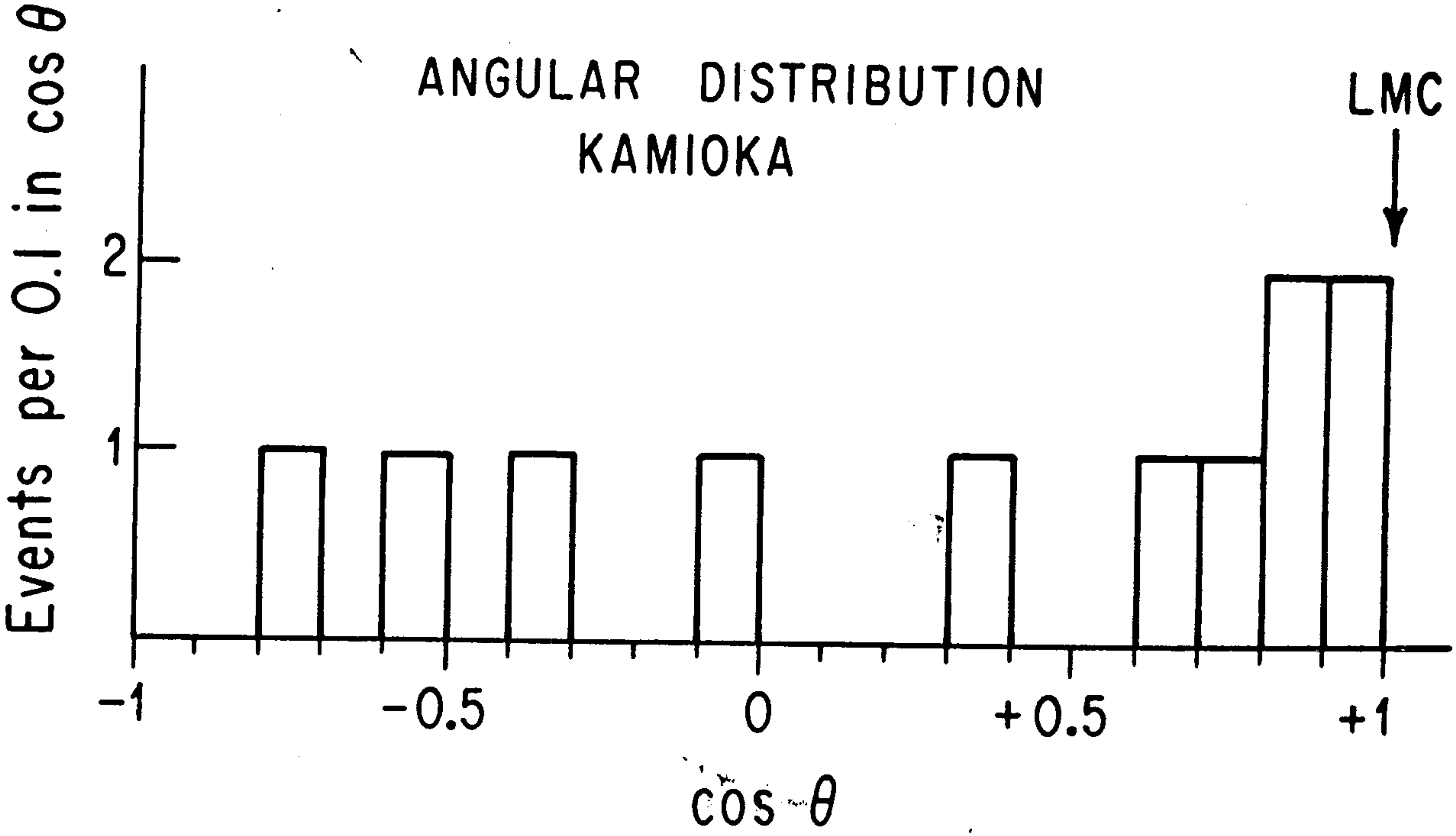

ANGULAR DISTRIBUTION
KAMIOKA
LMC
Events per 0.1 in cos θ
2
1
−1
−0.5
0
+0.5
+1
cos θ

FIGURE 3

Table 1: Neutrino Data		
Time (UT) February	Detector (threshold*/size)	# of Events (E-range/Duration)
23 2h 52m	Mt. Blanc (7 MeV/90 T)$^+$	5 (6–10 MeV/7 sec)
“”±1 min	Kamioka (8 MeV/2.14 kT)	1 (7 MeV/10 sec) (consistent w/background)
“”	IMB (30 MeV/5 kT)	none reported
“”	Baksan (11 MeV/130 T)$^+$	none reported
23 7h 35m (± min)	Kamioka (7 MeV/90 T)	11 (7–35 MeV/13 sec)
23 7h 35m	IMB (30 MeV/5 kT)	8 (20–40 MeV/4 sec)
“”	Baksan (11 MeV/130 T)$^+$	3 (12–17 MeV/10 sec)***
“”	Mt. Blanc (7 MeV/90 T)$^+$	2 (7–9 MeV/13 sec)
sum of pulses	Homestake ν_e (0.7 MeV/615 T)**	0
	Optical	
23 9h 25m	lack of sighting	$m_\nu \gtrsim 8$ magnitude
23 10h 40m	photograph	$m_\nu = 6$ magnitude
24 10h 53m	discovery	$m_\nu = 4.8$ magnitude

*Threshold is when efficiency drops to $\lesssim 50\%$ (sub-threshold events are therefore possible).

$^+$These detectors are liquid scintalators with $H_{2n+n}C_n$, thus have $\sim$ 1.39 more free protons than H_2O detectors of same mass.

**The Homestake detector is only sensitive to ν_e's. It is made of C_2Cl_4.

***Three fiducial volume events, 5 total volume events.

Table 2: Analysis of Mt. Blanc Event

	Mean E (MeV)	$\langle\sigma\rangle$	eff. $T_{\bar{\nu}_e}$ (MeV)	$\epsilon_{\bar{\nu}_e}(\times 10^{52}\text{ergs})$	$\epsilon_{total}(10^{53}\text{ergs})$
all 5 events					
with 5 MeV cut-off	8.4	10.2	1.6	96 ± 42	64 ± 29
with 7 MeV cut-off	8.4	10.2	0.3	$8 \pm 3 \times 10^{11}$	$5 \pm 2 \times 10^{11}$
3 high events					
with 7 MeV cut-off	9.3	11.1	0.9	$3.6 \pm 1.6 \times 10^{3}$	$2.4 \pm 1 \times 10^{3}$
with 5 MeV cut-off	9.3	11.1	1.8	44 ± 19	30 ± 13