Modern Physics Letters A, Vol. 5, No. 31 (1990) 2581–2592

WHAT HAVE WE LEARNED FROM SN 1987A?

GEORG G. RAFFELT
Max-Planck-Institut für Physik, Postfach 401212, 8000 München 40, Germany

Received 10 April 1990
Revised 2 August 1990

We review the constraints on the properties of neutrinos, axions, majorons, light supersymmetric particles, and other constraints on fundamental physics that have been derived on the basis of the neutrino observations of SN 1987A.

1. Introduction

About three years after the neutrino pulse from the celebrated supernova 1987A was observed in the IMB[1] and Kamiokande II[2] water Čerenkov detectors, extracting information relevant for particle physics from these measurements is still an ongoing effort. Nevertheless, it seems to be a good time to review what has been learned from these measurements. While the number of papers written on this subject is large, the actual number of different arguments used to derive particle physics constraints is quite limited so that it is most economical to proceed in terms of these arguments, reviewing the applications on specific cases for each method.

The first and simplest set of arguments is based on the fact the $\bar{\nu}_e$ pulse and perhaps the de-leptonization pulse were observed, constraining various mechanisms that could have removed neutrinos from the pulse such as decays. Equally, the non-observation of a γ-ray burst in coincidence with the neutrino burst constrains radiative decays. More intricate arguments involve signal dispersion, either between the photon- and neutrino-observations, between the $\bar{\nu}_e$ and the (questionable) ν_e observations, or the intrinsic dispersion of the $\bar{\nu}_e$ burst, constraining various effects that could cause signal dispersion such as a non-zero neutrino mass. Finally, the inferred cooling time scale of the newborn neutron star precludes effective cooling by other agents and thus yields constraints on the emission of new particles from the supernova core such as axions. We now review these arguments in order.

2. Absence of Non-Standard Signatures

2.1. *Observation of the $\bar{\nu}_e$-pulse*

The neutrino burst from a supernova is expected to consist of two major phases: the initial de-leptonization burst which contains a large fraction of ν_e, and the cooling phase where (anti-) neutrinos of all flavors are about equally emitted. The water Čerenkov detectors would register the de-leptonization burst by virtue of the reaction $\nu_e + e^- \rightarrow \nu_e + e^-$ where the scattered electron is strongly forward peaked, while the cooling signal is registered by $\overline{\nu}_e + p \rightarrow n + e^+$ with an essentially isotropic positron signal. Most if not all of the detected neutrinos must have been $\overline{\nu}_e$'s, in agreement with theoretical expectations, while the first event in Kamiokande II is consistent with being due to a ν_e.

The observation of the $\overline{\nu}_e$'s precludes that these particles have decayed on their way from the supernova to us, yielding a constraint on their lifetime of[3]

$$\tau_{\nu_e}/m_{\nu_e} \gtrsim 6 \times 10^5\,\mathrm{s/eV}\,. \qquad (1)$$

However, even this simple result must be interpreted with care. In the presence of neutrino mixing, and as neutrino flavors are expected to mix if neutrinos have masses, only the heavy ν_e admixtures would decay and could well violate this bound.[3]

Also, the $\overline{\nu}_e$'s were not removed from the burst by excessive scattering on cosmic background neutrinos, majorons, dark matter particles, etc., leading to constraints on "secret interactions" of neutrinos.[4] Specifically, $g_x \lesssim 10^{-3}$ is a constraint on the neutrino-majoron Yukawa coupling.

The number of detected events and their energies correspond well to what is expected theoretically so that no significant number of $\overline{\nu}_e$'s can have been *added* to the pulse. There are several mechanisms that could contribute to the pulse apart from the direct emission at the source. Other neutrino flavors ($\overline{\nu}_\mu$, $\overline{\nu}_\tau$) could decay in flight, producing $\overline{\nu}_e$ as decay products.[5] Equally, resonant oscillations in the supernova mantle could transform other flavors into detectable $\overline{\nu}_e$.[6] More interestingly, neutrinos with anomalous magnetic or electric dipole or transition moments would suffer helicity flips in the SN core because of their interactions with charged particles. If the flipped states are non-interacting, implying that neutrinos are of the Dirac type, they would freely escape from the core as opposed to the left-handed states which are trapped and are thus radiated from the neutrino sphere. The right-handed states could oscillate back into interacting, left-handed neutrinos in the galactic magnetic field and would contribute to the detected signal. Because the energies of these events would be characteristic of the core temperature rather than the temperature at the neutrino sphere, the absence of such high energy $\overline{\nu}_e$'s allows one to derive a bound on the electromagnetic moment of[7,8]

$$\mu_\nu \lesssim 10^{-12}\,\mu_\mathrm{B}\,, \qquad (2)$$

where $\mu_\mathrm{B} = e/2m_e$ is the Bohr-magneton. This results relies on the assumption that neutrinos are Dirac particles, and that the Earth is not coincidentally at a node of the neutrino magnetic oscillation pattern.

Other particles besides neutrinos which could have been emitted from the SN could also contribute to the signal. This fact, for example, yields constraints on axion properties.[9]

The angular distribution of the $\overline{\nu}_e$ signal is consistent with isotropy only at a confidence level of a few percent. An alternate interpretation in terms of hypothetical X^0 particles which would produce a forward biased signal[10] probably runs afoul of a constraint based on the helium burning lifetime of globular cluster stars.[11]

2.2. *Possible observation of the ν_e-pulse*

The high energy and forward direction of the first Kamiokande II event is consistent with it due to the de-leptonization ν_e-pulse although this interpretation cannot be made statistically significant. Nevertheless, it is interesting to discuss what could be learned if this interpretation were correct because a future galactic SN observation may well provide a significant signal. Many authors have shown that the large range of length scales and density conditions in the SN mantle allows resonant neutrino oscillations to occur for a large range of neutrino mass differences and mixing angles.[6,12–23] Specifically, if the ν_e's survive in the SN, there is very little, if any, parameter space where the MSW effect could solve the solar neutrino puzzle. Moreover, if μ or τ neutrinos were the dark matter of the universe, they would have a mass of about (50 – 100) eV. The observation of the ν_e-pulse would then require the ν_e-ν_μ or ν_e-ν_τ mixing angle to be constrained by $\sin^2 2\theta \lesssim 10^{-7}$, i.e., this observation would practically rule out neutrino dark matter.[14]

The unambiguous identification of the de-leptonization pulse would also allow one to determine whether a SN consists of matter or anti-matter.[24]

2.3. *Non-observation of a γ-ray burst*

During the time of the neutrino observations, the solar maximum mission (SMM) satellite was operational and registered a normal background flux of γ-rays.[25] The absence of a γ burst in association with the neutrino burst allows one to constrain radiative neutrino decays[25–29]:

$$\tau_{\nu_e}/m_{\nu_e} \gtrsim 2\times 10^{15}\,\mathrm{s/eV}\,, \tag{3}$$

where τ is the *radiative* lifetime only. A similar bound pertains to μ and τ neutrinos if they are lighter than about 20 eV, while for larger masses the bound is less restrictive because the photon spectrum would be spread out in time. A general bound, valid for all families, is[26]

$$\frac{\tau_\nu}{1\ \mathrm{sec}} \gtrsim \begin{cases} 2\times 10^{15}(m_\nu/1\ \mathrm{eV}) & \text{if} \quad m_\nu \lesssim 20\ \mathrm{eV}\,, \\ 3\times 10^{16} & \text{if} \quad 20\ \mathrm{eV} \lesssim m_\nu \lesssim 100\ \mathrm{eV}\,, \\ 8\times 10^{17}(1\ \mathrm{eV}/m_\nu) & \text{if} \quad 100\ \mathrm{eV} \lesssim m_\nu \lesssim 1\ \mathrm{MeV}\,. \end{cases} \tag{4}$$

More interestingly, these results can be expressed in terms of an electric or magnetic transition moment, μ_ν. Because the radiative decay width is given by $\tau_\nu^{-1} = \mu_\nu^2 m_\nu^3/8\pi$, assuming that the final-state neutrino is much lighter than the initial state, we find

$$\frac{\mu_\nu}{\mu_\mathrm{B}} \lesssim \begin{cases} 1\times 10^{-8}(1\ \mathrm{eV}/m_\nu)^2 & \text{if} \quad m_\nu \lesssim 20\ \mathrm{eV}\,, \\ 5\times 10^{-10}(1\ \mathrm{eV}/m_\nu) & \text{if} \quad m_\nu \gtrsim 100\ \mathrm{eV}\,, \end{cases} \tag{5}$$

where $\mu_B = e/2m_e$ is the Bohr magneton.

3. Dispersion Effects

3.1. Photons vs. anti-neutrinos

The optical sighting of SN 1987A followed the detection of the $\bar{\nu}_e$ burst by only a few hours, a delay which is expected on the basis of the simple reasoning that some time must pass before the mantle of a supernova "notices" the collapse of the inner core. Hence the two signals must have propagated through space with an almost identical velocity, i.e., the speed of light and that of neutrinos are equal to within[30,31]

$$\left|\frac{c_\nu - c_\gamma}{c_\gamma}\right| \lesssim 2 \times 10^{-9} , \tag{6}$$

assuming an uncertainty of ± 3 h in the relative duration of the transit times from the Large Magellanic Cloud to us. This was interpreted as the most stringent test of special relativity to date in the sense that it proves with high precision the universality of a relativistic limiting velocity.

Moreover, this result can also be interpreted as testing the weak equivalence principle of general relativity.[32] In the post Newtonian approximation, one predicts that a gravitational potential $U(\mathbf{r})$ delays a light signal by an amount

$$\Delta t = -2 \int_E^A U[\mathbf{r}(t)]\, dt , \tag{7}$$

where the integral is taken along the trajectory $\mathbf{r}(t)$ of the beam between the points of emission (E) and absorption (A). This delay is the same for neutrinos and photons to within

$$\left|\frac{\Delta t_\nu - \Delta t_\gamma}{\Delta t_\gamma}\right| < (0.7 - 4) \times 10^{-3} , \tag{8}$$

where the uncertainty reflects the uncertain modelling of the gravitational potential between Earth and SN 1987A. This result has been used to constrain the parameters of a specific model of C and P violating gravitational forces,[33] and to constrain the parameters of a class of non-metric theories of gravity.[34]

3.2. Neutrinos vs. anti-neutrinos

Assuming that the de-leptonization pulse was observed, one may also constrain the difference in transit time between ν_e and $\bar{\nu}_e$ and thus confirm the equivalence principle between matter and anti-matter.[35] However, such results cannot be made statistically significant on the basis of the SN 1987A observations.

3.3. Intrinsic dispersion of the $\bar{\nu}_e$-pulse

The previous arguments constrained differences in the transit time of different particle species where the transit time was assumed to be energy-independent, an assumption which is valid for massless particles if general relativity is correct. However, the most likely effect of signal propagation over large distances is dispersion due to an energy-dependent speed of propagation. The most widely discussed effect of this sort is that of a non-zero mass for the electron neutrino.[36–48] The main problem in extracting information about the dispersion of the signal is the unknown behavior of the source which must be modelled according to some theoretical assumptions. The most complete discussion is probably that of Loredo and Lamb[48] who, employing a consistent statistical methodology, included the background of the detectors in their analysis and tested a variety of emission models. They favor an exponential cooling model with a fixed radius of the neutrino sphere and a temperature which decreases as $T_{\nu\text{-sphere}} = T_0 e^{-t/4\tau}$ so that τ is a time scale for the decrease of the neutrino luminosity. Moreover, they used the parameter $\alpha \equiv (R_{\text{obs}}/10\text{ km})\,(50\text{ kpc}/D)g^{1/2}$, where D is the distance to SN 1987A and g is a weight factor which is unity if only left-handed, massless neutrinos of any given flavor are being emitted (three flavors were assumed to exist). They took the mass of the electron neutrino as a free parameter in order to allow for signal dispersion, and they introduced two separate offset times for the IMB and Kamiokande II detectors between the arrival of the first neutrinos and the first detected event. Hence they allowed the following six parameters to vary in order to achieve a maximum likelihood result: T_0, τ, α, m_{ν_e}, t_{off} (IMB), and t_{off} (KII). The best-fit values are given in the first column of Table 1, where we also show the inferred values for the neutron star proper radius, R, the total amount of binding energy, E_b, and the number of expected neutrino detections, N_{det}, in each detector. The neutrino mass is found to be limited by

$$m_{\nu_e} < 23\text{ eV} \quad (95\%\text{ C. L.})\,, \tag{9}$$

Table 1. Maximum likelihood results inferred from the observed neutrino pulse of SN 1987A, using an exponential cooling model, and including detector background events in the analysis (after Loredo and Lamb[48]). If the neutrino mass is taken to be a free parameter, the best-fit result is $m_{\nu_e} = 0$, with all other parameters having the best-fit values shown in the first column. In the second column, we show the results if m_{ν_e} is assumed to be 23 eV, a value which is the upper limit at the 95% C. L.

m_{ν_e}	0	23 eV
T_0	4.47 MeV	4.84 MeV
τ	4.15 sec	2.96 sec
α	2.26	2.06
t_{off} (KII)	0	3.57 sec
t_{off} (IMB)	0	0.85 sec
Fitted parameters above, Inferred parameters below.		
E_b	2.86×10^{53}erg	2.33×10^{53}erg
R	22.6 km	20.6 km
N_{det} (KII)	12.5	11.5
N_{det} (IMB)	5.51	6.14

and, taking this value as a fixed choice, the best-fit values for the remaining 5 parameters are given in the second column of Table 1.

The absence of any anomalous dispersion of the neutrino pulse can be used to constrain other neutrino properties. A small electric charge, e_ν, would bend the neutrino path in the galactic magnetic field, leading to an energy-dependent time delay:

$$\frac{\Delta t}{t} = \frac{e_\nu^2 (B_\perp d_{\rm coh})^2}{6E_\nu^2}, \tag{10}$$

where $B_\perp$ is the transverse magnetic field and $d_{\rm coh}$ its coherence length. This leads to a constraint of[49]

$$\frac{e_\nu}{e} \lesssim 3 \times 10^{-17} \left(\frac{1\,\mu{\rm G}}{B_\perp}\right)\left(\frac{1\,{\rm kpc}}{d_{\rm coh}}\right). \tag{11}$$

Speculating further one may consider some sort of "fifth force charge", q_ν, for neutrinos. If electrons, protons, or dark matter particles also carry such a charge, $q_{e,p,x}$, the bending of the neutrino trajectory in the fifth force field of the galaxy would lead to an energy-dependent time delay, yielding a constraint[35,50,51]

$$|q_{e,p,x}\, q_\nu| \lesssim 3 \times 10^{-40}. \tag{12}$$

4. Duration of Neutrino Emission

4.1. *The general argument*

The most intricate way to constrain particle properties arises from the observed duration of the neutrino pulse. While the first burst of neutrinos is emitted when the shock front passes through the outer core, dissociating heavy nuclei and thereby unlocking neutrinos which where trapped by coherent interactions, the long tail of the neutrino signal is associated with cooling, i.e., emission from the neutrino sphere which is powered by thermal energy stored in the inner core. The existence of a direct cooling channel of this inner region, such as the emission of axions or right-handed neutrinos, would deprive the late cooling phase of energy and thus would shorten the observed neutrino pulse.

The general argument is best illustrated considering axion emission. Axions are pseudoscalar particles which, for the purpose of this argument, are taken to interact with neutrons and protons with a common Yukawa coupling strength, g_a, which is the only free parameter of the problem. For very small values of g_a, axions will play no role; but with increasing coupling strength, their emission from the inner core by bremsstrahlung processes, $N + N \rightarrow N + N + a$, will begin to compete with neutrino cooling. Of course, once g_a exceeds some critical value, axions will be trapped and emitted from an "axio-sphere", the radius of which increases with increasing g_a. Eventually, axions will be trapped so effectively that, again, their contribution to the cooling of the SN core is negligible and the neutrino signal assumes its standard duration. This general behavior[52–56] is shown in Fig. 1 on the basis of detailed numerical investigations.[57,58]

From Fig. 1 it is quite apparent that a large range of g_a values can be excluded on the basis of the observed duration of the neutrino signal which agrees well with theoretical expectations. It must be stressed that here we are considering the time scale of neutrino *emission* at the source, while the detectors register a pulse which conceivably could have been lengthened by dispersion effects. In this regard it is remarkable that even assuming m_{ν_e} as large as 23 eV, i.e., its 95% C. L. upper limit, does not substantially reduce the inferred duration of neutrino emission at the source – see Table 1. Hence we can be reasonably sure that neutrinos were indeed emitted over several seconds.

While it is clear that a large range of g_a values can be excluded, no statistical analysis has been performed that would allow one to state a confidence level for a certain range of excluded parameters. Hence, the ultimately stated bounds are somewhat arbitrary.

Another caveat applies to the "trapping regime" of the new particles. In this case one may expect that the new particles may play an important role during the infall phase and shock formation of a SN collapse, an issue that was addressed only by a small number of authors in the context of the majoron bounds[59] and bounds on neutrino dipole moments.[7] Hence it is not obvious that parameters allowed by the cooling argument on the trapping side would remain allowed if one took account of these effects. Conversely, one may speculate that some sort of new particle physics may actually help to couple the observed SN explosions to the gravitational collapse, the "supernova problem" which is notoriously difficult to resolve using standard physics.

We may now go through a list of several cases to which variations of this general cooling argument have been applied. However except for the axion case where detailed numerical investigations are available, all other results are based on simple analytic criteria such as demanding that the luminosity in new

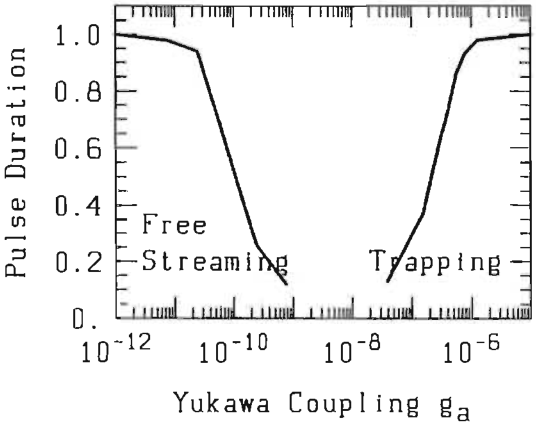

Fig. 1. Duration of the neutrino pulse from a supernova, taking axion emission into account.[57,58] The curve is an average of the results relevant for the IMB and Kamiokande detectors. For a very small or very large Yukawa coupling to protons and neutrons, g_a, the duration is normalized to unity, reflecting the standard value where axion cooling is irrelevant. In the intermediate regime between free streaming and trapping no numerical results are available.

particles should not exceed a certain plausible value, typically 2×10^{53} erg/s at typical conditions in the SN core. It is impossible to assess, in general, how such criteria compare with the requirement that the neutrino pulse was not shortened by more than a specified factor.

4.2. *Axions*

We have already used axions as an example to illustrate the general argument above – they were the first case to which this reasoning was applied.[52–58] On the basis of Fig. 1 and requiring that the neutrino pulse was not shortened by more than a factor of two, one finds an excluded regime of

$$10^{-10} \lesssim g_a \lesssim 10^{-7} \, . \tag{13}$$

This result, however, is afflicted with rather large uncertainties because of the uncertain calculation of the emission rates in a dense nuclear medium, but a detailed discussion of these uncertainties is beyond the scope of this brief review. While they are difficult to quantify, they probably do not exceed a factor of a few on either side of the excluded range.

The axion Yukawa couplings to protons and neutrons are, in general,

$$g_{an} = \frac{c_n m_\mathrm{N}}{f_a} \, , \quad g_{ap} = \frac{c_p m_\mathrm{N}}{f_a} \, , \tag{14}$$

where m_N is the nucleon mass, f_a the axion decay constant, and c_n and c_p are model-dependent numbers of order unity. The axion mass is given by

$$\frac{m_a}{1\ \mathrm{eV}} = \frac{0.60 \times 10^7\ \mathrm{GeV}}{f_a} \, . \tag{15}$$

Hence one may express the bounds on the axion couplings as a bound on the axion mass if one specifies an axion model and thus the numbers c_n and c_p. For hadronic axions, the excluded regime is approximately

$$3 \times 10^{-3}\,\mathrm{eV} \lesssim m_a \lesssim 3\ \mathrm{eV} \, . \tag{16}$$

The axion bounds as well as other cases are afflicted with substantial uncertainties in the calculation of the emission rates which are largely due to the difficulties in dealing with many-body effects in a dense nuclear medium. One such effect is the possibility of a phase transition to a quark-gluon plasma so that quarks and gluons, rather than nucleons and mesons, would have to be considered the fundamental excitations of the medium. The axion emission rates from a quark-gluon plasma are currently being investigated by two groups of authors.[60–62]

4.3. *Right-handed neutrinos*

Another class of particles that could drain the SN core of energy are right-handed (RH) neutrinos, i.e., non-interacting states. These particles could be an entire new class of sterile neutrinos, particularly if the known neutrinos are of the Majorana type. If the known neutrinos are Dirac particles, they could simply be the helicity-flipped states.

The simplest way to produce RH neutrinos is by the same processes which produce LH states, assuming there exist RH weak interactions on some level. Assuming further that these RH interactions have the same structure as the LH interactions, one may easily derive bounds on the RH Fermi-constant, G_{RH}. On the basis of the modified urca processes, $n + n \rightarrow n + p + e^- + \overline{\nu}_{e\mathrm{R}}$ and $e^- + n + p \rightarrow n + n + \nu_{e\mathrm{R}}$, which involve charged currents, one finds in the free streaming regime[53]

$$G_{\mathrm{RH}} \lesssim 3 \times 10^{-5} G_{\mathrm{F}} \,. \tag{17}$$

The trapping regime is of much less interest because it overlaps with a regime excluded by laboratory data. On the basis of another emission process, $e^- + p \rightarrow n + \nu_{\mathrm{R}}$, one finds a similar constraint.[63] In the standard left-right symmetric models, this result can be translated into a bound on the mass of RH gauge bosons, $m_{W_{\mathrm{R}}}$, and the standard W_{R}-W_{L}-mixing angle, ζ,[63]

$$[\zeta^2 + (m_{W_{\mathrm{L}}}/m_{W_{\mathrm{R}}})^4]^{1/2} \lesssim 3 \times 10^{-5} \,. \tag{18}$$

Similarly, one may constrain RH neutral currents on the basis of bremsstrahlung processes, $N + N \rightarrow N + N + \overline{\nu}_{\mathrm{R}} \nu_{\mathrm{R}}$, yielding[53] $G_{\mathrm{RH}} \lesssim 10^{-4}$, although a somewhat weaker bound was reported by other authors.[63,64] Moreover, in standard left-right symmetric models, the RH neutral current has vector structure and thus does not contribute to nucleon bremsstrahlung,[64] leaving us with a much less efficient emission process, $e^+e^- \rightarrow \overline{\nu}_{\mathrm{R}} \nu_{\mathrm{R}}$. For E_6 models, where RH neutrino masses can be expected to be naturally small, the neutral-current bremsstrahlung rates are not suppressed, and a detailed analysis and interpretation of the constraints is available.[64,65] The constraint on the neutrino charge radius,[66] we believe, should be discussed in a unified picture with RH neutral current interactions since in the framework of electroweak gauge theories a neutrino charge radius is a problematic concept.

In the previous cases one had to assume that the mass of the RH neutrinos was small enough for them to be thermally emitted from the SN. The following arguments rely on various processes of flipping the helicity of Dirac neutrinos, thereby transforming an interacting (LH) state into a sterile (RH) state of the same mass. For ν_e and ν_μ the following constraints are thus valid without restriction, while for ν_τ with a laboratory limit on its mass of 35 MeV it is most likely that a mass range exists near this limit at which the following bounds can be evaded.

The simplest way to flip the helicity is by a mass term, i.e., the neutrinos trapped in the SN core will develop RH components if they have a Dirac mass. The resulting neutrino luminosity will be so large that one can infer a bound[53, 67–69]

$$m_\nu \lesssim 20 \text{ keV}. \tag{19}$$

It was claimed that the excluded mass range reaches up to ~ 35 MeV, with a substantial uncertainty, however, so that τ neutrinos with masses near their laboratory limit are still allowed by this argument.

If Dirac neutrino masses are induced by the interaction with a light Higgs field, neutrino scattering by means of virtual Higgs exchange would flip helicities. Hence one can exclude a certain region in the plane of Higgs vs. neutrino masses.[70]

If neutrinos had a magnetic or electric dipole moment, interactions with charged particles in the SN core would also flip the helicity, yielding a bound[7,8,71,72] of

$$\mu_\nu \lesssim 10^{-12}\mu_B\,, \tag{20}$$

where $\mu_B = e/2m_e$ is the Bohr magneton.

Finally, the helicity flip in the gravitational field of the nascent neutron star in the context of novel gravitational interactions was also discussed, allowing one to constrain the parameters of such models.[73]

4.4. *Other particles*

In supersymmetric models with light photinos, these particles would be emitted by nucleon bremsstrahlung processes. The cross-section is $\propto m_{\tilde{q}}^{-4}$, leading to a constraint on the squark mass of[64,74,75]

$$m_{\tilde{q}} \gtrsim 1\ \text{TeV}\,. \tag{21}$$

Many authors,[59,76–86] have discussed the effect of majorons on supernovae, although most of them concentrated on the triplet majoron model which is now excluded on the basis of the measured Z^0 width. The most recent investigation,[86] however, concerns a detailed account of bounds on the singlet majoron model, excluding a large range of neutrino masses and vacuum expectation values.

5. Summary

We have summarized a large number of constraints on particle physics parameters that have been derived on the basis of the observed neutrino pulse from SN 1987A. The most interesting results, in our opinion, are those which address well-motivated particles or particle properties, and which cannot be evaded easily. In this regard we favor the axion bounds, and the bounds on neutrino masses, both from dispersion and cooling arguments. However, apart from such more personal choices it is clear that SN 1987A has given us a large number of interesting results to be refined by observation of neutrinos from a galactic supernova which we are anxiously awaiting!

Acknowledgments

The author thanks B. Sadoulet for soliciting this review as a contribution to the DPF 90 conference, Houston, Texas, January, 1990, and for the support granted by the Center for Particle Astrophysics, Berkeley. Moreover, this research was supported by grants from NASA, DOE, and the Institute for Geophysics and Planetary Physics (Livermore).

References

1. R. M. Bionta *et al.*, *Phys. Rev. Lett.* **58** (1987) 1494. C. B. Bratton *et al.*, *Phys. Rev.* **D37** (1988) 3361.
2. K. S. Hirata *et al.*, *Phys. Rev. Lett.* **58** (1987) 1490; *Phys. Rev.* **D38** (1988) 448.
3. J. A. Frieman, H. E. Haber, and K. Freese, *Phys. Lett.* **200B** (1988) 115.
4. E. W. Kolb and M. S. Turner, *Phys. Rev.* **D36** (1987) 2895.
5. J. M. Soares and L. Wolfenstein, *Phys. Rev.* **D40** (1989) 3666.
6. L. Wolfenstein, *Phys. Lett.* **194B** (1987) 197.
7. D. Nötzold, *Phys. Rev.* **D38** (1988) 1658.
8. R. Barbieri and R. N. Mohapatra, *Phys. Rev. Lett.* **61** (1988) 27.
9. D. Seckel *et al.*, Report, 1990.
10. J. C. van der Velde, *Phys. Rev.* **D39** (1989) 1492.
11. G. Raffelt, *Phys. Rev.* **D38** (1988) 3811.
12. S. P. Mikheev and A. Yu. Smirnov, *Sov. Phys. JETP* **64** (1986) 4 [*Zh. Eksp. Teor. Fiz.* **91** (1986) 7].
13. P. O. Lagage, M. Cribier, J. Rich, and D.Vignaud, *Phys. Lett.* **193B** (1987) 127.
14. J. Arafune, M. Fukugita, T. Yanagida, and M. Yoshimura, *Phys. Lett.* **194B** (1987) 477.
15. J. Arafune and M. Fukugita, *Phys. Rev. Lett.* **59** (1987) 367.
16. J. Arafune, M. Fukugita, T. Yanagida, and M.Yoshimura, *Phys. Rev. Lett.* **59** (1987) 1864.
17. T. P. Walker and D. N. Schramm, *Phys. Lett.* **195B** (1987) 331.
18. D. Nötzold, *Phys. Lett.* **196B** (1987) 315.
19. T. K. Kuo and J. Pantaleone, *Phys. Rev.* **D37** (1988) 298.
20. H. Minakata, H. Nunokawa, K. Shiraishi, and H. Suzuki, *Mod. Phys. Lett.* **A2** (1987) 827.
21. H. Minakata and H. Nunokawa, *Phys. Rev.* **D38** (1988) 3605.
22. H. Minakata and H. Nunokawa, TMUP-HEL-8913 (1989).
23. S. P. Rosen, *Phys. Rev.* **D37** (1988) 1682.
24. A. V. Barnes, T. J. Weiler, and S. Pakvasa, *Astrophys. J.* **323** (1987) L31.
25. E. L. Chupp, W. T. Vestrand, and C. Reppin, *Phys. Rev. Lett.* **62** (1989) 505.
26. E. W. Kolb and M. S. Turner, *Phys. Rev. Lett.* **62** (1989) 509.
27. F. von Feilitzsch and L. Oberauer, *Phys. Lett.* **200B** (1988) 580.
28. M. Takahara and K. Sato, *Mod. Phys. Lett.* **A2** (1987) 293.
29. T. Hatsuda, C. S. Lim, and M. Yoshimura, *Phys. Lett.* **203B** (1988) 462.
30. M. J. Longo, *Phys. Rev.* **D36** (1987) 3276.
31. L. Stodolsky, *Phys. Lett.* **201B** (1988) 353.
32. L. Krauss and S. Tremaine, *Phys. Rev. Lett.* **60** (1988) 176.
33. L. D. Almeida, G. E. A. Matsas, and A. A. Natale, *Phys. Rev.* **D39** (1989) 677.
34. A. A. Coley and S. Tremaine, *Phys. Rev.* **D38** (1988) 2927.
35. S. Pakvasa, W. A. Simmons, and T. J. Weiler, *Phys. Rev.* **D39** (1989) 1761.
36. L. F. Abbott, A. de Rújula, and T. P. Walker, *Nucl. Phys.* **B299** (1988) 734.
37. E. N. Adams, *Phys. Rev.* **D37** (1988) 2047.
38. W. D. Arnett and J. L. Rosner, *Phys. Rev. Lett.* **58** (1987) 1906.
39. J. N. Bahcall and S. L. Glashow, *Nature* **326** (1987) 476.
40. A. Burrows and J. M. Lattimer, *Astrophys. J.* **318** (1987) L63.
41. A. Burrows, *Astrophys. J.* **328** (1988) L51.
42. H. -Y. Chiu, K. L. Chan, and Y. Kondo, *Astrophys. J.* **329** (1988) 326.
43. R. Cowsik, *Phys. Rev.* **D37** (1988) 1685.
44. E. W. Kolb, D. L. Tubbs, and D. A. Dicus, *Astrophys. J.* **255** (1982) L57.
45. S. Midorikawa, H. Terazawa, and K. Akama, *Mod. Phys. Lett.* **A2** (1987) 561.
46. K. Sato and H. Suzuki, *Phys. Rev. Lett.* **58** (1987) 2722; *Phys. Lett.* **196B** (1987) 267.
47. D. N. Spergel and J. N. Bahcall. *Phys. Lett.* **200B** (1988) 366.
48. T. J. Loredo and D. Q. Lamb, *Ann. N. Y. Acad. Sci.* **571** (1989) 601.

49. G. Barbiellini and G. Cocconi, *Nature* **329** (1987) 21.
50. J. A. Grifols, E. Massó, and S. Peris, *Phys. Lett.* **207B** (1988) 493.
51. G. Fiorentini and G. Mezzorani, Report, 1988.
52. J. Ellis and K. A. Olive, *Phys. Lett.* **193B** (1987) 525.
53. G. Raffelt and D. Seckel, *Phys. Rev. Lett.* **60** (1988) 1793.
54. R. Mayle, J. R. Wilson, J. Ellis, K. Olive, D. N. Schramm, and G. Steigman, *Phys. Lett.* **203B** (1988) 188.
55. R. Mayle, J. R. Wilson, J. Ellis, K. Olive, D. N. Schramm, and G. Steigman, *Phys. Lett.* **219B** (1989) 515.
56. M. S. Turner, *Phys. Rev. Lett.* **60** (1988) 1797.
57. A. Burrows, M. S. Turner, and R. P. Brinkmann, *Phys. Rev.* **D39** (1989) 1020.
58. A. Burrows, T. Ressell, and M. S. Turner, Report, 1990, to be published in *Phys. Rev.* **D**.
59. G. Fuller, R. Mayle, and J. R. Wilson, *Astrophys. J.* **332** (1988) 826.
60. T. Altherr, Report HU-TFT-90-3 (January, 1990), to be published in *Z. Phys.* **C**.
61. T. Altherr, Report HU-TFT-90-36 (June, 1990).
62. J. Ellis and P. Salati, CERN-TH-5693 (1990).
63. R. Barbieri and R. N. Mohapatra, *Phys. Rev.* **D39** (1989) 1229.
64. J. A. Grifols and E. Massó, *Nucl. Phys.* **B331** (1990) 244.
65. J. A. Grifols, E. Massó, and T. G. Rizzo, Report MAD/PH/549 (December, 1989).
66. J. A. Grifols and E. Massó, *Phys. Rev.* **D40** (1989) 3819.
67. K. J. F. Gaemers, R. Gandhi, and J. M. Lattimer, *Phys. Rev.* **D40** (1989) 309.
68. A. Perez and R. Gandhi, *Phys. Rev.* **D41** (1990) 2374.
69. J. A. Grifols and E. Massó, CERN-TH. 5632/90 (January, 1990).
70. J. A. Grifols, Report UAB-FT-239 (March, 1990).
71. J. M. Lattimer and J. Cooperstein, *Phys. Rev. Lett.* **61** (1988) 23.
72. I. Goldman, Y. Aharonov, G. Alexander, and S. Nussinov, *Phys. Rev. Lett.* **60** (1988) 1789.
73. D. Choudhury, N. D. Hari Dass, and M. V. N. Murthy, *Class. Quantum Grav.* **6** (1989) L167.
74. J. Ellis, K. A. Olive, S. Sarkar, and D. W. Sciama, *Phys. Lett.* **215B** (1988) 404.
75. J. A. Grifols, E. Massó, and S. Peris, *Phys. Lett.* **220B** (1989) 591.
76. H. M. Georgi, S. L. Glashow, and S. Nussinov, *Nucl. Phys.* **B193** (1981) 297.
77. G. B. Gelmini, S. Nussinov, and M. Roncadelli, *Nucl. Phys.* **B209** (1982) 157.
78. D. A. Dicus, E. W. Kolb, and D. L. Tubbs, *Nucl. Phys.* **B223** (1983) 532.
79. Y. Aharonov, F. T. Avignone III, and S. Nussinov, *Phys. Lett.* **200B** (1988) 122.
80. Y. Aharonov, F. T. Avignone III, and S. Nussinov, *Phys. Lett.* **D37** (1988) 1360.
81. Y. Aharonov, F. T. Avignone III, and S. Nussinov, *Phys. Lett.* **D39** (1989) 985.
82. R. V. Konoplich and M. Yu. Khlopov, *Sov. J. Nucl. Phys.* **47** (1988) 565 [*Yad. Fiz.* **47** (1986) 891].
83. J. A. Grifols, E. Massó, and S. Peris, *Phys. Lett.* **215B** (1988) 593.
84. Z. G. Berezhiani and A. Yu. Smirnov, *Phys. Lett.* **220B** (1989) 279.
85. K. Choi, C. W. Kim, J. Kim, and W. P. Lam, *Phys. Rev.* **D37** (1988) 3225.
86. K. Choi and A. Santamaria, MPI-PAE/PTh 75/89, to be published in *Phys. Rev.* **D**.