Mass bounds on new neutral gauge bosons in E_6 models from SN 1987A data

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E_6 models contain right-handed neutrinos which, if light, should be emitted during stellar collapse. Their interactions with matter involve the exchange of an extra neutral gauge boson Z' . We use data from SN 1987A to place bounds on the mass of this new gauge boson which in E_6 models can be parametrized in terms of a mixing angle θ . The limits obtained are, in general, much stronger than those from accelerator data and lie, except for a narrow band around $\theta = -14.48^\circ$, at the several TeV mass scale.

Superstring theory has prompted, on the phenomenological front, an exhaustive study of E_6 models.¹ The chiral matter in E_6 models is contained in a 27-dimensional representation which implies the existence of new degrees of freedom added to the 15 conventional fields of the standard model (SM). In particular, a right-handed (RH) neutrino is included in this menagerie of new particles. This extra neutrino degree of freedom, if light enough, could affect the energetics of a supernova (SN) collapse. Gravitational collapse involves the release of about 10^{53} ergs.² In February 1987 terrestrial detectors³ were struck by neutrinos originating in a SN explosion in the Large Magellanic Cloud (LMC). The energy inferred from these neutrino events exceeded 2×10^{53} ergs, emitted over a period of 1–10 sec. In models of stellar collapse, this energy is less than 4×10^{53} ergs, and thus there is some room left for other sources of energy emission. As can be easily deduced from the above-mentioned figures, any exotic drain of energy should have a rate smaller than 2×10^{53} ergs/sec. This fact can be used to constrain extra exotic sources of energy depletion.

In a previous paper⁴ (paper I), two of the present authors studied the emission of ordinary left-handed (LH) neutrinos, RH neutrinos, and photinos. The techniques set up to describe light-particle emission included the computation of the so-called Rosseland mean free path which is an energy-averaged inverse opacity that enters the theta-type equation describing the energy diffusion process.⁵ The observational data on SN 1987A were then correctly reproduced by blackbody radiation from the surface of last neutrino scattering, the so-called neutrinosphere, determined by means of the Rosseland parameter. Of course, particles other than ordinary neutrinos, e.g., RH neutrinos or photinos, may interact so weakly with matter in the stellar core that they may freely stream out of the star instead of diffusing through the very dense

stellar matter. Diffusion and free streaming are two extreme regimes which a light particle produced in the core can experience. It is the coupling of the particle to matter that determines which regime applies.

The detailed consideration of both qualitatively different regimes, diffusion and free streaming, for the (light) photinos of supersymmetry (SUSY) and (light) RH neutrinos of left-right-symmetric and E_6 models led to relevant upper and lower bounds on squark and slepton masses and on the mass of new Z' neutral gauge bosons.⁵ In particular, when E_6 models were considered, three particular models were chosen to show the techniques and to stress the power of the methods used to obtain strong constraints on the extra Z' neutral gauge boson. In this paper we would like to complete this previous work on E_6 neutral-gauge-boson mass limits. We remind the reader that our bounds apply if either ν_R is a light Majorana neutrino or if it forms a Dirac fermion when combined with ν_L .

Let us parametrize the lowest-lying extra neutral boson as an arbitrary mixture of Z_ψ and Z_χ , Z_ψ being the result of E_6 breaking down to $SO(10) \times U(1)_\psi$ and Z_χ raising from the breakdown to $SU(5) \times U(1)_\chi$. Then

$$Z' = Z_\psi \cos\theta - Z_\chi \sin\theta. \quad (1)$$

Particular models are found by giving specific values to the mixing angle θ . Here we keep θ arbitrary for the sake of generality.

The production of ordinary neutrino pairs in the stellar core proceeds mainly through nucleon-nucleon bremsstrahlung ($NN \rightarrow NN\nu\bar{\nu}$). RH neutrinos will be produced via this very same mechanism, but with a Z' boson mediating the four-fermion neutral-current interaction instead of the usual SM Z boson. We start considering the case where the neutrino mean free path is larger than the di-

mensions of the star. The ν_R luminosity can be calculated following the procedure discussed in paper I. Essentially, it involves the evaluation of the energy loss via $NN \rightarrow NN \nu_R \bar{\nu}_R$, where the nucleons are to be thermalized in the nondegenerate approximation and in the non-relativistic (NR) limit. In this NR limit only the axial-vector couplings to nucleons are operative and we find, for the energy/volume time output,

$$q_{NN}(\nu_R) = 48g^4 \tan^4 \theta_W g_{\nu_R}^2 g_A^2 X_n \times [A_n^2 X_n + 6(A_n^2 + A_p^2) X_p] \frac{1}{M_{Z'}^4} \dot{\epsilon}_{NN}, \quad (2)$$

where

$$\dot{\epsilon}_{NN} = \frac{16}{1155\pi^{11/2}} \rho^2 f^4 \left[\frac{m}{m_\pi} \right]^4 \frac{T^{11/2}}{m^{9/2}}$$

(ρ is the matter density, f is the strong pion-nucleon coupling constant, m is the effective nucleon mass, and T is the temperature of the inner stellar core), the nucleon axial-vector couplings are

$$A_n = - \left[\frac{5}{2} \right]^{1/2} c_\theta + \frac{1}{\sqrt{6}} s_\theta, \\ A_p = \left[\frac{5}{2} \right]^{1/2} c_\theta + \frac{1}{\sqrt{6}} s_\theta,$$

the RH-neutrino coupling is

$$g_{\nu_R} = \frac{1}{6} \left[\frac{5}{2} \right]^{1/2} c_\theta + \frac{5}{2\sqrt{6}} s_\theta,$$

$g_A = 1.259$, and $X_{n,p}$ are the mass fractions of neutrons (protons) in the core of the star.

In a simple model⁶ with an inner core of radius $R = 10$ km, constant density $\rho = 8 \times 10^{14}$ g/cm³, and temperature $T = 60$ MeV, the ν_R luminosity (energy per unit time) is

$$Q_{\nu_R} = \frac{4}{3} \pi R^3 q_{NN}(\nu_R). \quad (3)$$

The power output resulting from Eq. (3) cannot be arbitrarily large. As stated before, there could be some 2×10^{53} ergs extra energy available which is depleted over a period that lasts for several seconds. We will hence take as a nominal upper limit for the ν_R luminosity the value 2×10^{53} ergs/sec. We feel that this is a very safe number since, should the ν_R luminosity be anywhere close to this value, the core would cool so quickly that the duration of the observed pulse would be of order 1 sec or less in contradiction with observation. Indeed, detailed numerical studies for the similar problem of axion cooling⁷ have shown that a main consequence of the back reaction of axion emission is the shortening of the time pattern of the neutrino burst. The neutrino emission times are drastically reduced even with a modest axion energy drain. We expect also in our case that even for ν_R luminosities much lower than 2×10^{53} ergs/sec, the feedback on ordinary neutrino emission would produce a timing for the $\bar{\nu}$ signal substantially different from what was observed.

The condition $Q_{\nu_R} \leq 2 \times 10^{53}$ ergs/sec enables us to put a lower limit on $M_{Z'}$ appearing in Eq. (2). This limit is shown in Fig. 1 as a function of the mixing angle θ . This is a reliable lower bound because we took into account the uncertainties associated with the core parameters and displayed in Fig. 1 the most conservative result. In fact, the curve shown is the result of dividing by a factor 2.4 the limits obtained using the core parameters given before Eq. (3). This factor is the uncertainty associated with the Z' mass which was estimated in paper I.⁴ These results should be compared with the lower limits on Z' from the Collider Detector at Fermilab⁸ (CDF). Their data imply that, for the whole range of the angle θ , the mass of the extra Z' should be larger than about 300 GeV. Hence, except for a very narrow band around $\theta = -14.48^\circ$, the bounds from SN 1987A are much stronger than the ones obtained from accelerator data.

The bounds discussed so far correspond to the free-streaming regime because Z' masses are so large that matter is transparent to RH neutrinos. Indeed, their mean free path in matter is easily seen to be much larger than the core radius R for all values of θ when the Z' mass values in Fig. 1 are used. We should now consider the other possibility, namely, that the mass of the Z' boson is so low that RH neutrinos are forced to interact substantially with the matter in the core. RH neutrinos in this case slowly diffuse as they random walk to the RH neutrinosphere from where they are thermally emitted. Consideration of diffusion leads to an upper bound for $M_{Z'}$. Indeed, the mass of the Z' gauge boson should be small enough so that RH neutrinos are emitted from a sufficiently large neutrinosphere where the temperature is less and hence the luminosity is lower. If $M_{Z'}$ exceeds a given critical value [see Eq. (13) below], the luminosity goes beyond the permitted 2×10^{53} ergs/sec. One should keep in mind, however, that the diffusion approximation breaks down for sufficiently small coupling strengths

FIG. 1. Lower bound on $M_{Z'}$ as a function of θ obtained from RH-neutrino emission via nucleon-nucleon bremsstrahlung.

since in this case the neutrinos will escape freely no matter how light the Z' boson is.

To compute the RH-neutrino luminosity in this regime, we need to know the radius of the ν_R -sphere and its temperature, and then apply the Stefan-Boltzmann law to this radiating sphere. To do so, a temperature profile for the star is required and we adopt the profile used in paper I,

$$T(r) = \begin{cases} 15 \text{ MeV } \rho_{14}^{2/3}, & \rho_{14} \geq 2, \\ 20 \text{ MeV } \rho_{14}^{1/3}, & \rho_{14} \leq 2, \end{cases} \quad (4)$$

and first introduced by Turner in Ref. 6. Here $\rho_{14} = (r/1.5 \times 10^6 \text{ cm})^{-7}$ is the density in units of 10^{14} g/cm^3 .

As mentioned before, one can find the neutrinosphere R_{ν_R} from an energy-averaged mean free path $\bar{\lambda}(r)$:

$$\bar{\lambda}(r) \equiv \int dE (\sigma n)^{-1} E^4 \frac{e^{E/T(r)}}{(e^{E/T(r)} + 1)^2} \bigg/ \int dE E^4 \frac{e^{E/T(r)}}{(e^{E/T(r)} + 1)^2}, \quad (5)$$

where σ is the neutrino-nucleon cross section and n is the density of nucleon targets in the core. Indeed, from the calculation of the optical depth

$$\tau(r) = - \int_{\infty}^r \frac{dr}{\bar{\lambda}(r)}, \quad (6)$$

and the condition

$$\tau(r = R_{\nu_R}) = \frac{2}{3}, \quad (7)$$

we obtain the neutrinosphere R_{ν_R} . The $\nu_R N \rightarrow \nu_R N$ cross section is given by

$$\sigma(\nu_R N \rightarrow \nu_R N) = \frac{25}{288} g^4 \tan^4 \theta_W \omega(\theta, X_n, X_p) \frac{E^2}{M_{Z'}^4}, \quad (8)$$

with

$$\omega(\theta, X_n, X_p) \equiv \left[\frac{1}{2\sqrt{6}} c_\theta + \frac{5}{2\sqrt{10}} s_\theta \right]^2 \left\{ \left[\frac{8}{5} s_\theta^2 + 3g_A^2 \left[-\frac{3}{\sqrt{6}} c_\theta + \frac{1}{\sqrt{10}} s_\theta \right]^2 \right] X_n + \left[\frac{2}{5} s_\theta^2 + 3g_A^2 \left[\frac{3}{\sqrt{6}} c_\theta + \frac{1}{\sqrt{10}} s_\theta \right]^2 \right] X_p \right\}.$$

We now introduce Eq. (6) into Eq. (5) to obtain $\bar{\lambda}(r)$ explicitly,

$$\bar{\lambda}(r) = 1.6 \times 10^{-5} \frac{[M_{Z'} (\text{GeV})]^4}{\omega(\theta, X_n, X_p)} \left[\frac{r}{1.5 \times 10^6 \text{ cm}} \right]^{35/3}, \quad (9)$$

from which we get

$$\frac{R_{\nu_R}}{1.5 \times 10^6 \text{ cm}} = 2.3 \omega(\theta, X_n, X_p)^{3/2} [M_{Z'} (\text{GeV})]^{-3/8}. \quad (10)$$

The last step is to calculate the luminosity Q_{ν_R} from

$$Q_{\nu_R} = \frac{7\pi^3}{240} T_{\nu_R}^4 R_{\nu_R}^2, \quad (11)$$

using both the neutrinosphere radius R_{ν_R} and the temperature profile Eq. (4) for $T_{\nu_R} = T(R_{\nu_R})$.

This Q_{ν_R} is bounded by

$$Q_{\nu_R} \leq 2 \times 10^{53} \text{ ergs/sec}, \quad (12)$$

to be consistent with observation.

This allows us to conclude

$$M_{Z'} \leq 80 \text{ GeV } \omega^{1/4}(\theta, X_n, X_p). \quad (13)$$

We note that the dependence of $M_{Z'}$ on the assumed luminosity bound is mild ($M_{Z'} \simeq Q_{\nu_R}^{4/11}$). We show the

FIG. 2. Upper bound on $M_{Z'}$ as a function of θ obtained from thermal radiation of RH neutrinos.

FIG. 3. Lower bound on $M_{Z'}$ as a function of θ obtained from S_L^c emission via nucleon-nucleon bremsstrahlung.

FIG. 4. Lower bound on $M_{Z'}$ as a function of θ obtained from S_L^c emission via e^+e^- annihilation.

limit implied by Eq. (13) in Fig. 2. Two remarks are in order. First, for θ close to -14.48° (where RH-neutrino couplings to Z' vanish), the results are nonsensical since, in this case, as already stressed before, the diffusion approximation fails. Actually, there are no limits in this case because both the opacity and production cross section go to zero. Second, accelerator data already exclude the allowed band of Z' masses. As a consequence, we are left with the lower bounds shown in Fig. 1.

The limits derived in this paper are valid in the case when the RH neutrino has a mass much smaller than the temperature of the core. This is probably not a too contrived restriction since Majorana masses are difficult to implement in an E_6 scenario and Dirac neutrino masses emerge naturally from an E_6 -invariant superpotential.

For the $SO(10)$ -singlet S_L^c in the 27 representation of E_6 , the lightness assumption is possibly far less justified, but if one assumes that S_L^c is also very light, then one can

apply the same techniques and obtain bounds on Z' from S_L^c emission in SN 1987A collapse. In addition, if S_L^c and ν_L form a Dirac state, these same limits will apply. Figure 3 shows the bounds on $M_{Z'}$ from nucleon-nucleon bremsstrahlung, and Fig. 4 displays the bounds on the mass of the Z' which follow from e^+e^- annihilation via Z' exchange.

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