Energizing gamma ray bursts via $Z'$ mediated neutrino heating

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Abstract

The energy deposition in stellar explosions due to the pair annihilation of neutrinos ($\nu\bar{\nu} \rightarrow e^+e^-$) can energize events such as Type II supernovae, merging neutron stars, gamma ray bursts (GRBs) etc. This neutrino heating can be further enhanced by modifying the background geometry over that of the Newtonian spacetime. However, even then the observed energy in GRBs cannot be achieved. In this paper, we explore if the inclusion of the contribution due to an extra $Z'$ gauge boson in the neutrino pair annihilation process can explain the energy required for a GRB. We compute the expression for the energy deposition coming from the $Z'$ mediated process in Newtonian, Schwarzschild, and Hartle-Thorne spacetime backgrounds. We show that the contribution due to the $Z'$ mediated process can enhance the energy deposition rate significantly even in the Newtonian background. The inclusion of the general relativistic corrections and rotation in the background spacetime near neutron stars can further enhance the energy deposition rates. From the observed energy of GRBs, we obtain constraints on the mass and the gauge coupling of the extra $Z'$ for the three background spacetimes mentioned.
I. INTRODUCTION

It has long been realized that the neutrino pair annihilation process plays a significant role in depositing energy to violent stellar processes such as type II supernovae [1–3], merging neutron stars [4], binary neutron stars in the last stable orbit [5], Gamma Ray Bursts (GRBs) [6–20], etc. It is believed that most of the energy in such stellar explosions is carried away by the three flavors of neutrinos. The emission of a huge number of neutrinos with luminosity $L_\nu \sim 10^{52}\text{erg/s}$ makes the stellar objects cool. A fraction of such huge neutrino flux can also deposit energy into the stellar envelope through neutrino pair annihilation ($\nu_i \bar{\nu}_i \rightarrow e^+ e^-, i = e, \mu, \tau$), neutrino lepton scattering, and neutrino baryon capture [17]. This mechanism is termed as neutrino heating. The process $\nu_i \bar{\nu}_i \rightarrow e^+ e^-$ is important for collapsing neutron stars, binary neutron stars in their last stable orbit, and r process nucleosynthesis [21]. This annihilation process can also continuously give energy to the radiation bubble and is a possible source of powering gamma ray bursts (GRBs) [6]. GRBs are high energy explosions that have been observed from cosmological distances [13, 22–24]. These are the most energetic phenomenon represented by intense and prompt $\gamma$-ray emissions. The energy emission associated with the GRBs is $\sim 10^{52}\text{erg}$ [25]. However, in a Newtonian background the neutrino pair annihilation process cannot provide the observed energy.

The effect of including different background spacetimes near such strong gravity regimes has also been underscored in the literature. This is also apparent from the values of $\frac{2GM}{R}$ which for instance is $\sim 0.7$ for collapsing neutron stars, and $\sim 0.4$ for supernovae calculations and hence, one cannot neglect the effect of General Relativity (GR) in the strong gravity regime [5]. Here, $G$ denotes Newton’s gravitational constant, $M$ denotes the mass of the neutron star, and $R$ denotes the distance scale. Since, the neutron stars are rotating, the rotation parameter in the background metric also has to be included. It has been shown that in the Schwarzschild background, the neutrino heating is enhanced up to a factor 4 for type II supernova, and by up to a factor 30 for collapsing neutron stars relative to the Newtonian result [6]. The rotation on the other hand reduces the energy deposition by 38% compared to the non rotating case [10]. The energy deposition for $\nu \bar{\nu} \rightarrow e^+ e^-$ has also been studied in other alternative theories of gravity [17] which are well motivated to solve dark
matter, dark energy, cosmological singularity problems etc. The quintessence field can also enhance the energy deposition in these systems as discussed in [18]. The radial variation of the temperature for black hole accretion disk in modified gravity theories can enhance the energy deposition by one order magnitude with respect to GR as is recently discussed in [26]. The energy deposition can be enhanced as well due to the presence of topological defects such as global monopole [27]. All these studies have considered only the Standard Model (SM) contribution to the neutrino pair annihilation process.

Although the SM is very well established and can explain majority of the experimental results, there are several motivations of going beyond it. This includes neutrino mass, dark matter, matter-antimatter asymmetry of the universe, muon $g - 2$ etc. Hence, probing signatures and implications of Beyond Standard Model (BSM) physics in different observations and experiments is of cardinal importance.

In this paper, we consider for the first time, the possibility of additional contribution from BSM physics to neutrino pair annihilation process in the context of increasing the energy budget of GRBs. In particular we focus on the scenario where the SM is extended by a general $U(1)_X$ gauge group. Such extensions are well motivated from the viewpoint of generation neutrino masses, since the cancellation of gauged and mixed gauged gravity anomalies warrant the presence of three right handed neutrinos. Such a scenario contains an additional neutral gauge boson ($Z'$) associated with the $U(1)_X$ symmetry. The special case of $U(1)_X$ is the $U(1)_{B-L}$ model which has been extensively studied in the literature [28-30]. The phenomenology of extra $Z'$ in these models is particularly interesting and constraints have been obtained on the mass and the coupling strength of $Z'$ from several experiments. This includes electroweak precession data [31], collider searches [32-37], neutrino-electron scattering experiments [38,39], beam dump experiments [40-42], SN1987A [43-46] etc. We note that the extra $Z'$ gauge boson present in such $U(1)_X$ models can enhance the cross section of the process $\nu\bar{\nu} \rightarrow e^+e^-$. This leads to an increased neutrino heating. We investigate the extent of the enhancement due to this additional contribution for Newtonian (Newt) as well as Schwarzschild (Sch) and Hartle-Thorne (HT) geometries. We also study the efficiency of the combined effects of modified background spacetime and BSM physics to explain the observed energy in GRBs. From the reported energy of GRBs ($10^{52}$ erg) [25,47],
we obtain constraints on the mass and the coupling of the \( Z' \) boson for the three background spacetime geometries mentioned above.

The paper is organized as follows. In Section II we obtain the total cross section and the energy deposition rate due to \( Z' \) mediated process in neutrino pair annihilation. In Section III we compute the angular integration in Newtonian, Schwarzschild, and Hartle-Thorne spacetime background for the neutrino heating process. The contribution of \( Z' \) mediated process to the neutrino heating mechanism in different spacetime backgrounds is calculated in Section IV. In Section V we present the constraints on \( Z' \) from neutrino heating in GRBs for the three spacetime backgrounds mentioned above. In Section VI our results are compared with the constraints on the mass and coupling of \( Z' \) coming from other experiments. Finally in Section VII we conclude and discuss our results.

In the following, we have used natural units \((c = 1, \hbar = 1)\), and \(G = 1\) throughout the paper.

II. NEUTRINO HEATING THROUGH \( Z' \)

We extend the SM gauge group \((SU(3)_c \times SU(2)_L \times U(1)_Y)\) by an additional \(U(1)_X\) gauge symmetry. Such a scenario includes an extra neutral gauge boson \((Z')\) associated with the \(U(1)_X\) symmetry. The latter is broken by an extra singlet scalar field \(\Phi\) and the gauge boson acquires mass. Three right handed neutrinos are needed in the model to cancel the gauge and gauge-gravity anomalies. The \(U(1)_X\) charges of the quarks and the leptons can be expressed in terms of that of the \(\Phi\) and the SM Higgs \((H)\). The charges for the scalars are chosen as \(2x_\Phi\) and \(\frac{x_\Phi}{2}\). The corresponding \(U(1)_X\) charge of lepton doublet is \(Q^l_X = (-\frac{1}{2}x_H - x_\Phi)\) and the charges of right handed electron and neutrino are \(Q^e_R = (-x_H - x_\Phi)\) and \(Q^N_R = -x_\Phi\) respectively. The particular choice of \(x_H = 0\) and \(x_\Phi = 1\) leads to the \(U(1)_{B-L}\) model \([48-52]\).

The general interaction Lagrangian of \(Z'\) gauge bosons with the leptons is

\[-\mathcal{L}_{int} \supset g' (Q^l_X \overline{L} \gamma^\mu I_L Z'_{\mu} + Q^e_R \overline{R} \gamma^\mu I_R Z'_{\mu}).\] (2.1)

The electron neutrino contributes to the neutrino annihilation process \(\nu_e \overline{\nu}_e \rightarrow e^+ e^-\) via charge current \((W)\), neutral current \((Z)\), and \(Z'\) mediated interactions whereas \(\nu_\mu\) and \(\nu_\tau\)
have only $Z$ and $Z'$ mediated interactions for the process $\nu_{\mu,\tau} \nu_{\mu,\tau} \rightarrow e^+e^-$. Since, all three flavors of neutrinos are there in a hot neutron star, they will all contribute to the energy deposition rate via the Feynman diagrams Fig. 1. The energy deposition rate per unit volume near a hot neutron star is given as \[ \dot{q}(r) = \int \int f_{\nu}(p_{\nu}, r)\sigma_{\nu} f_{\nu}(p_{\nu}, r) \left( E_{\nu} + E_{\nu} \right) \frac{d^3p_{\nu}}{E_{\nu} d^3p_{\nu}}. \] where $E_{\nu}$ denotes the energy of neutrino and $f_{\nu}$ corresponds to their thermal energy distribution function in the phase space which is of Fermi-Dirac type and it is \[ f_{\nu} = \frac{2}{(2\pi)^3} \left( e^{E_{\nu}/kT} + 1 \right)^{-1}. \] Here $k$ denotes the Boltzmann constant and $T$ denotes the neutrino temperature. The neutrino velocity is denoted as $v_{\nu}$ and $\sigma$ denotes the cross section in the rest frame. Eq. 2.2 is true for any flavor of neutrinos. Since, the term in the first bracket of Eq. 2.2 is a Lorentz invariant quantity, we can calculate its value in the centre of mass frame for $\nu_e \bar{\nu}_e \rightarrow e^+e^-$ process. In the $U(1)_X$ model, it can be expressed as

\begin{align*}
(\sigma|v_{\nu_{\mu}} - v_{\nu_{\tau}}|E_{\nu_{\mu}} E_{\nu_{\tau}})_{U(1)_X} &= \left[ \frac{G_F^2}{3\pi} \left( 1 + 4 \sin^2 \theta_W + 8 \sin^4 \theta_W \right) + \frac{4g^4}{6\pi M_{Z'}^2} \left\{ \frac{3}{4} x_H + x_{\phi} \right\}^2 \left( \frac{x_H}{4} \right)^2 \right]
+ \frac{4G_F g^2}{3\sqrt{2}\pi M_{Z'}^2} \left( x_{\phi} + \frac{x_H}{4} \right)^2 \left[ \left( \frac{3}{4} x_H + x_{\phi} \right) \left( -\frac{1}{2} + 2 \sin^2 \theta_W \right) + \frac{x_H}{8} \right]
+ \frac{4G_F g^2}{3\sqrt{2}\pi M_{Z'}^2} \left( x_{\phi} + \frac{x_H}{2} \right)^2 \left( E_{\nu_{\mu}} E_{\nu_{\tau}} - p_{\nu_{\mu}} \cdot p_{\nu_{\tau}} \right)^2,
\end{align*}

(2.3)
where for the region of interest we neglect the mass of the electron, since, the energy of neutrino is greater than 10 MeV; \( G_F = 1.166 \times 10^{-5} \) GeV\(^{-2} \) is the Fermi constant and the \( \theta_W \) is the Weinberg angle whose value is \( \sin^2 \theta_W = 0.23 \) \([\text{53}]\). The first term in Eq. 2.3 corresponds to the \( W \) and \( Z \) mediated SM processes, the second term is due to the \( Z' \) contribution only. The third term arises because of the interference between \( Z \) and \( Z' \) mediated diagrams, while the fourth term stems from the interference between the \( W \) and \( Z' \) mediated processes. For muon and tau type of neutrinos, only \( Z \) and \( Z' \) mediated processes will contribute. Hence, for the scattering \( \nu_{\mu,\tau} \rightarrow e^+e^- \), we obtain

\[
(\sigma | V_{\nu_{\mu,\tau}} - V_{\nu_{\mu,\tau}} | E_{\nu_{\mu,\tau}} E_{\nu_{\mu,\tau}})_{U(1)_X} = \frac{G_F^2}{3\pi} (1 - 4 \sin^2 \theta_W + 8 \sin^4 \theta_W) + \frac{4g'^4}{6\pi M_{Z'}^4} \left\{ \left( \frac{3}{4} x_H + x_\Phi \right)^2 + \left( \frac{x_H}{4} \right)^2 \right\} \times \frac{4G_F g'^2}{3\sqrt{2\pi} M_{Z'}} \left[ \left( \frac{3}{4} x_H + x_\Phi \right) \left( -\frac{1}{2} + 2 \sin^2 \theta_W \right) + \frac{x_H}{8} \right] \times \left( E_{\nu_{\mu,\tau}} E_{\nu_{\mu,\tau}} - P_{\nu_{\mu,\tau}} P_{\nu_{\mu,\tau}} \right)^2.
\]

(2.4)

In what follows, we focus on the \( U(1)_{B-L} \) model which has fixed values of \( x_H \) and \( x_\Phi \) as \( x_H = 0 \) and \( x_\Phi = 1 \). The results can easily be generalized to the \( U(1)_X \) case by using suitable values of \( x_H \) and \( x_\Phi \).

Putting \( P_\nu = E_\nu \Omega_\nu \), and \( d^3p_\nu = E_\nu^2 dE_\nu d\Omega_\nu \) in the direction of \( \Omega_\nu \), where \( d\Omega_\nu \) denotes the solid angle, and assuming \( T_\nu = T_\tau = T \), we find

\[
\int \int f_\nu f_\tau (E_\nu + E_\tau) E_\nu^3 E_\tau^3 dE_\nu dE_\tau = \frac{21}{2(2\pi)^6} \pi^4 (kT)^9 \zeta(5).
\]

(2.5)

Hence, from Eq. 2.2 the energy deposition rate due to the electron neutrino pair annihilation process becomes

\[
\dot{q}_{\nu_e}(r) = \frac{21}{2(2\pi)^6} \pi^4 (kT_{\nu_e}(r))^9 \zeta(5) \times \left[ \frac{G_F^2}{3\pi} (1 + 4 \sin^2 \theta_W + 8 \sin^4 \theta_W) + \frac{4g'^4}{6\pi M_{Z'}^4} \right] \times \frac{4G_F g'^2}{3\sqrt{2\pi} M_{Z'}^2} \left[ -\frac{1}{2} + 2 \sin^2 \theta_W \right] \Theta_{\nu_e}(r),
\]

(2.6)

where the angular integration \( \Theta(r) \) is defined as

\[
\Theta(r) = \int \int (1 - \Omega_\nu \cdot \Omega_\tau)^2 d\Omega_\nu d\Omega_\tau.
\]

(2.7)
Similarly the energy deposition due to muon and tau type neutrino annihilations is given by

\[ \dot{q}_{\nu_{\mu,\tau}}(r) = \frac{21}{2(2\pi)^6} \pi^4 (kT_{\nu_{\mu,\tau}}(r))^9 \zeta(5) \times \left[ \frac{G_F^2}{3\pi} (1 - 4 \sin^2 \theta_W + 8 \sin^4 \theta_W) + \frac{4g'^4}{6\pi M_{Z'}^4} + \frac{4G_F g'^2}{3\sqrt{2}\pi M_{Z'}^2} \left( -\frac{1}{2} + 2 \sin^2 \theta_W \right) \right] \Theta_{\nu_{\mu,\tau}}(r). \] (2.8)

In the SM limit \( \frac{g'}{M_{Z'}} \to 0 \), we get back the earlier result \[20, 54\]

\[ \dot{q}(r) = \frac{7G_F^2 \pi^3 \zeta(5)}{2(2\pi)^6} (kT)^9 \Theta(r)(1 \pm 4 \sin^2 \theta_W + 8 \sin^4 \theta_W), \] (2.9)

where the + sign is for \( \nu_e \bar{\nu}_e \) pair and the minus sign is for \( \nu_{\mu} \bar{\nu}_{\mu} \) and \( \nu_{\tau} \bar{\nu}_{\tau} \) pairs.

### III. NEUTRINO HEATING IN DIFFERENT SPACETIME BACKGROUND

The geodesic outside a slowly rotating neutron star with only a dipole correction on a static star is governed by the HT metric \[55\]

\[ ds^2 = -\left(1 - \frac{2M}{r} + \frac{2J^2}{r^4}\right)dt^2 + \left(1 - \frac{2M}{r} + \frac{2J^2}{r^4}\right)^{-1} dr^2 + r^2 d\theta^2 + \left(d\phi - \frac{2J}{r^3} dt\right)^2, \] (3.1)

where \( r \) denotes the distance from the origin, \( \phi \) is the longitude, \( M \) is the mass of the neutron star, and \( J \) is the specific angular momentum. For \( J = 0 \), Eq. 3.1 reduces to the Schwarzschild metric and for both \( J = M = 0 \), we have the flat or the Newtonian metric. Since, we have considered the planar motion for the massless particle so we take \( \theta = \frac{\pi}{2} \) and the null geodesic \( g_{\mu\nu}V^\mu V^\nu = 0 \), where \( V^\mu = \frac{dX^\mu}{dt} \) and \( X^\mu \equiv (t, r, \theta, \phi) \). Hence, the geodesic equation of the HT metric for the planar motion becomes

\[ -\left(1 - \frac{2M}{r} - \frac{2J^2}{r^4}\right) \dot{t}^2 + \left(1 - \frac{2M}{r} + \frac{2J^2}{r^4}\right)^{-1} \dot{r}^2 + r^2 \dot{\phi}^2 - \frac{4J r}{r^3} \dot{t} \dot{\phi} = 0. \] (3.2)

We can also derive the generalized momenta as

\[ p_t = -\left(1 - \frac{2M}{r} - \frac{2J^2}{r^4}\right) \dot{t} - \frac{2J}{r} \dot{\phi} = -E, \quad p_r = \left(1 - \frac{2M}{r} + \frac{2J^2}{r^4}\right)^{-1} \dot{r}, \quad p_\phi = -\frac{2J \dot{t}}{r} + r^2 \dot{\phi} = L, \] (3.3)

where \( E \) denotes the energy per unit mass of the system and \( L \) denotes the angular momentum per unit mass. We can solve Eq 3.3 and obtain the expressions of \( \dot{t} \) and \( \dot{\phi} \) in terms of
\( L \) and \( E \) as

\[
i = \left( E - \frac{2JL}{r^3} \right) \left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^{-1}, \quad \dot{\phi} = \left[ \frac{L}{r^2} \left( 1 - \frac{2M}{r} - \frac{2J^2}{r^4} \right) + \frac{2JE}{r^3} \right] \left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^{-1}.
\] (3.4)

Using Eq. 3.3 we can write Eq. 3.2 as

\[- E \dot{t} + L \dot{\phi} + \left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^{-1} \dot{r}^2 = 0. \] (3.5)

Putting the expressions of \( \dot{t} \) and \( \dot{\phi} \) from Eq 3.4, we can write Eq. 3.5 as

\[
\frac{1}{r^4} \left( \frac{dr}{d\phi} \right)^2 = \frac{\left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^2}{\left( 1 - \frac{2M}{r} - \frac{2J^2}{r^4} + \frac{2JE}{r^3} \right)^2} \left[ \frac{E^2}{L^2} - \frac{4JE}{Lr^3} - \frac{1}{r} \left( 1 - \frac{2M}{r} - \frac{2J^2}{r^4} \right) \right].
\] (3.6)

The local Lorentz tetrad for the given metric can be written as [10]

\[
e_i^\mu = \begin{bmatrix}
\left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^{\frac{3}{2}} & 0 & 0 & 0 \\
0 & \left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^{-\frac{3}{2}} & 0 & 0 \\
0 & 0 & r & 0 \\
-\frac{2J}{r^2} & 0 & 0 & r \sin \theta
\end{bmatrix},
\] (3.7)

where the upper index denotes the row and the lower index denotes the column. The tangent of the angle \( (\theta_r) \) between the trajectory and the tangent vector is the ratio between the radial and longitudinal velocity and can be expressed in terms of \( \frac{dr}{d\phi} \) as

\[
\left( \frac{dr}{d\phi} \right)^2 = \left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^{\frac{3}{2}} \left[ \left( 1 - \frac{2M}{r} - \frac{2J^2}{r^4} \right)^{-1} \right] \tan \theta_r.
\] (3.8)

Comparing Eq. 3.6 and Eq. 3.8 we obtain the impact parameter \( b = \frac{L}{E} \) as

\[
b = \left[ \frac{2J}{r^3} + \left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^{\frac{3}{2}} \right]^{-1}.
\] (3.9)

If the neutrino is emitted tangentially \( (\theta_R = 0) \) from the neutrinosphere of radius \( R_{\nu_i} \), then its trajectory defined by an angle \( \theta_r \) with radius \( r \) is given by

\[
\cos \theta_r^{\nu_i} = \frac{R_{\nu_i}^3 r^2 \left( 1 - \frac{2M}{r} + \frac{2J^2}{r^4} \right)^{\frac{3}{2}}}{2J (r^3 - R_{\nu_i}^3) + R_{\nu_i}^2 r^3 \left( 1 - \frac{2M}{R_{\nu_i}} + \frac{2J^2}{R_{\nu_i}^4} \right)^{\frac{3}{2}}}.
\] (3.10)
where $i = e, \mu, \tau$. In the Newtonian limit, Eq. 3.10 becomes $(\cos \theta^\nu_r)_{\text{Newt}} = \frac{R_{\nu_i}}{r}$, and in the non rotating Schwarzschild limit Eq. 3.10 becomes $(\cos \theta^\nu_r)_{\text{Sch}} = \frac{R_{\nu_i}}{r} \sqrt{\frac{1 - 2M}{1 - \frac{2M}{R_{\nu_i}}}}$.

Putting the expressions of $\dot{t}$ and $\dot{\phi}$ (Eq. 3.4) in Eq. 3.5 we obtain

$$\frac{E^2}{2} = \frac{r^2}{2} + V_{\text{eff}},$$

(3.11)

where the effective potential $V_{\text{eff}}$ defining the trajectory of a massless particle for a slowly rotating neutron star system is given as

$$V_{\text{eff}} = \frac{L^2}{2r^2} \left(1 - \frac{2M}{r} - \frac{2J^2}{r^4}\right) + \frac{2JL^2}{br^3}.$$  

(3.12)

To find the neutrinosphere radius which corresponds to the last stable circular orbit for the neutrinos, we have to impose $\frac{dV_{\text{eff}}}{dr} = 0$. Then from Eq. 3.12 we obtain

$$r^4 - r^3 \left(3M - \frac{6J}{b}\right) - 6J^2 = 0.$$  

(3.13)

If there is no rotation (Schwarzschild solution), then we get the neutrinosphere orbit at a radius $= 3M$. In that case, the massless neutrino which has a radius $< 3M$ is gravitationally bound. However, if we solve Eq. 3.13 for the rotating case, we can find at least one real root for which the neutrinosphere radius is $< 3M$. If the neutrinos are emitted tangentially from the neutrinosphere surface, then from Eq 3.9 we can also write,

$$r^6 - b^2r^4 + 2Mb^2r^3 + 2J^2b^2 - 4bJr^3 = 0.$$  

(3.14)

We can calculate the neutrinosphere radius and the impact parameter by solving Eq. 3.13 and Eq. 3.14 simultaneously for different values of $\frac{J}{M^2}$. The event horizon is calculated using $1 - \frac{2M}{R_{EH}} - \frac{2J^2}{R_{EH}^3} = 0$. We are interested in the domain $r > R_{\nu}$ for the energy deposition rate. In TABLE I we have summarized the event horizons, neutrinosphere radius and impact parameter for different values of $\frac{J}{M^2}$.

Now we compute the angular integration factor $\Theta(r)$ which appears in Eq. 2.7. Choosing $d\Omega = d\mu d\phi$, where $\mu = \sin \theta$ and $\Omega = (\mu, \sqrt{1 - \mu^2} \cos \phi, \sqrt{1 - \mu^2} \sin \phi)$, we can write

$$\Theta(r) = 4\pi^2 \int_0^1 \int_x^1 \left[1 - 2\mu \mu_\nu + \mu^2 \mu_\nu^2 + \frac{1}{2}(1 - \mu^2)(1 - \mu_\nu^2)\right] d\mu_\nu d\mu_\phi.$$  

(3.15)

Note that our expressions Eq.3.13 and Eq.3.14 are slightly different from those given in [10] due to some typographical errors.
| $\frac{J}{M^2}$ | $\frac{R_{EH}}{M}$ | $\frac{R_{\nu}}{M}$ | $\frac{b}{M}$ |
|--------------|------------------|------------------|----------|
| 0.1          | 2.002            | 2.882            | 4.990    |
| 0.2          | 2.009            | 2.759            | 4.770    |
| 0.3          | 2.022            | 2.632            | 4.534    |
| 0.4          | 2.038            | 2.500            | 4.278    |
| 0.5          | 2.057            | 2.363            | 3.998    |

**TABLE I**: Summary of the event horizon ($\frac{R_{EH}}{M}$), neutrinosphere radius ($\frac{R_{\nu}}{M}$), and impact parameter ($\frac{b}{M}$) for different values of $\frac{J}{M^2}$.

The solution of Eq. 3.15 is found as

$$\Theta(r) = \frac{2\pi^2}{3}(1 - x)^4(x^2 + 4x + 5),$$

where,

$$x = (\sin \theta^\nu_i)_{HT} = \left[1 - \frac{R^6_{\nu_i}r^4 \left(1 - \frac{2M}{r} + \frac{2J^2}{r^3}\right)}{2J(r^3 - R^3_{\nu_i}) + R^2_{\nu_i}r^3 \left(1 - \frac{2M}{R_{\nu_i}} + \frac{2J^2}{R^3_{\nu_i}}\right)^{\frac{1}{2}}}\right]^{\frac{1}{2}}.$$  \hspace{1cm} (3.17)

The temperature of the free streaming neutrinos at a radius $r$ is related to the temperature of the neutrino at the neutrinosphere by the gravitational redshift. The neutrino temperature varies linearly with the redshift as

$$T_{\nu_i}(r) = \sqrt{\frac{1 - \frac{2M}{R_{\nu_i}} - \frac{2J^2}{R^4_{\nu_i}}}{1 - \frac{2M}{r} - \frac{2J^2}{r^3}}} T_{\nu_i}(R_{\nu_i}),$$

whereas the observable quantity, the luminosity varies quadratically with the redshift as

$$L_{\text{obs}} = \left(1 - \frac{2M}{R_{\nu_i}} - \frac{2J^2}{R^4_{\nu_i}}\right) L_{\nu_i}(R_{\nu_i}).$$  \hspace{1cm} (3.19)
The neutrino luminosity for each $\nu_i$ species for a blackbody neutrino gas can also be written as

$$L_{\nu_i}(R_{\nu_i}) = 4\pi R_{\nu_i}^2 \frac{7}{16} a T_{\nu_i}^4(R_{\nu_i}),$$  \hspace{1cm} (3.20)$$

where $a = 0.663$ is the radiation constant in natural units.

Using Eq. 3.16, Eq. 3.18, Eq. 3.19, Eq. 3.20 we can calculate

$$T_{\nu_i}^0(r) \Theta_{\nu_i}(r) = \frac{\left(1 - \frac{2M}{R_{\nu_i}} - \frac{2J^2}{R_{\nu_i}^4}\right)^{\frac{3}{2}}}{\left(1 - \frac{2M}{r} - \frac{2J^2}{r^4}\right)^{\frac{3}{2}}} \left(\frac{7}{4} \pi a\right)^{-\frac{9}{4}} R_{\nu_i}^{\frac{9}{2}} L_{\text{obs}}^{\frac{9}{2}} (1 - x_{\nu_i})^4 (x_{\nu_i}^2 + 4x_{\nu_i} + 5).$$  \hspace{1cm} (3.21)$$

The energy deposition rate is enhanced with increasing $T_{\nu_i}^0(r) \Theta_{\nu_i}(r)$ which has distinct values in different background spacetimes. Note that, the limit $J = 0 \ (J = 0, \ M = 0)$ corresponds to Schwarzschild (Newtonian) background.

IV. CONTRIBUTION TO NEUTRINO HEATING FROM BOTH $Z'$ AND MODIFIED SPACETIME BACKGROUND

The total amount of energy deposition due to neutrino heating in the region beyond the neutrinosphere is given as

$$\dot{Q}_{\nu_i} = \int_{R_{\nu_i}}^{\infty} \dot{q}_{\nu_i} \frac{4\pi r^2 \, dr}{\sqrt{1 - \frac{2M}{r} - \frac{2J^2}{r^4}}}. \hspace{1cm} (4.1)$$

Using Eq. 2.6, Eq. 2.8, and Eq. 3.21, we can write the energy deposition rate for the electron neutrino in HT metric as

$$\dot{Q}_{\nu_e}^{\text{HT}} = \frac{28\pi^7}{(2\pi)^6} k^9(5) \times \left[ \frac{G_F^2}{3\pi} \left(1 + 4 \sin^2 \theta_W + 8 \sin^4 \theta_W\right) + \frac{4g'^4}{6\pi M_{Z'}^2} + \frac{4G_F g'^2}{3\sqrt{2} \pi M_{Z'}^2} \right] \left(1 - \frac{2M}{R_{\nu_e}} - \frac{2J^2}{R_{\nu_e}^4}\right)^{\frac{3}{2}} \left(\frac{7}{4} \pi a\right)^{-\frac{9}{4}} R_{\nu_e}^{\frac{9}{2}} L_{\text{obs}}^{\frac{9}{2}} \left(1 - x_{\nu_e}^{\text{HT}}\right)^4 (x_{\nu_e}^{\text{HT}}^2 + 4x_{\nu_e}^{\text{HT}} + 5). \hspace{1cm} (4.2)$$

$$\int_{1}^{\infty} \frac{y_{\nu_e}^2 \, dy_{\nu_e}}{(1 - \frac{2M}{y_{\nu_e} R_{\nu_e}} - \frac{2J^2}{(y_{\nu_e} R_{\nu_e})^4})^{\frac{3}{2}}} \left(1 - x_{\nu_e}^{\text{HT}}\right)^4 (x_{\nu_e}^{\text{HT}}^2 + 4x_{\nu_e}^{\text{HT}} + 5),$$
where \( y_{\nu_i} = \frac{r}{R_{\nu_i}} \) and \( x_{\nu_e}^{\text{HT}} \) is given by Eq. \( \text{3.17} \) for \( \nu_e \). Similarly, the energy deposition rates for \( \nu_\mu \) and \( \nu_\tau \) in the HT background become

\[
\dot{Q}_{\nu_\mu,\tau}^{\text{HT}} = \frac{28\pi^7}{(2\pi)^6} k^9 \zeta(5) \times \left[ \frac{G_F^2}{3\pi} (1 - 4\sin^2 \theta_W + 8\sin^4 \theta_W) + \frac{4g^\prime}{6\pi M^4_Z} + \frac{4G_F g^\prime}{3\sqrt{2}\pi M^2_Z} \right] \left( 1 - \frac{2M}{R_{\nu_\mu,\tau}} - \frac{2J^2}{R_{\nu_\mu,\tau}^4} \right)^2 \left( \frac{\pi a}{4} \right)^{-\frac{3}{2}} \frac{4 - 3}{4 R_{\nu_\mu,\tau}^3} L^9_{\text{obs}} \ 
\]

(4.3)

where \( x_{\nu_\mu,\tau}^{\text{HT}} \) is given by Eq. \( \text{3.17} \) for \( \nu_\mu,\tau \).

From Eq. \( \text{4.2} \) and Eq. \( \text{4.3} \) we can similarly obtain the energy deposition rates for three flavors of neutrinos in the Schwarzschild and the Newtonian background with proper choices of \( x_{\nu_e}^{\text{Sch}} \) and \( x_{\nu_e}^{\text{Newt}} \) by putting \( J = 0 \) and \( J = M = 0 \) in Eq. \( \text{3.17} \) respectively. The total energy deposition by all the neutrino species is \( \dot{Q}_{\nu} + \dot{Q}_{\nu_\mu,\tau} \) which can be calculated for all three types of spacetime mentioned above.

If we only include the W and Z mediated diagrams \( i.e \) only the SM contribution for the process \( \nu \bar{\nu} \rightarrow e^+ e^- \) and denote \( D = 1 \pm 4 \sin^2 \theta_w + 8 \sin^4 \theta_w \), where + sign is for \( \nu_e \bar{\nu}_e \) pair, and - sign is for \( \nu_\mu \bar{\nu}_\mu \) and \( \nu_\tau \bar{\nu}_\tau \) pairs then the total rate of energy deposition in HT spacetime background becomes

\[
\dot{Q}_{51} = 1.09 \times 10^{-5} F \left( \frac{M}{R}, \frac{J}{R^2} \right) D L_{51}^{9/4} R_6^{-3/2}, \ 
\]

(4.4)

where \( \dot{Q}_{51}^{\text{HT}} = \frac{\dot{Q}_{\nu_e}}{10^{44} \text{ erg/sec}}, \ L_{51} = \frac{L_{\text{obs}}}{10^{44} \text{ erg/sec}}, \ R_6 = \frac{R}{10 \text{ km}}, \) and the enhancement function

\[
F \left( \frac{M}{R}, \frac{J}{R^2} \right) = 3 \left( 1 - \frac{2M}{R} - \frac{2J^2}{R^4} \right)^{9/4} \int_1^\infty \frac{y^2 dy}{\left( 1 - \frac{2M}{yR} - \frac{2J^2}{yR^4} \right)^{5}} (1-x_{\text{HT}}^2) \left( x_{\text{HT}}^2 + 4x_{\text{HT}}^2 \right) + 5, \ 
\]

(4.5)

where we assume \( R_{\nu_e} \approx R_{\nu_\mu,\nu_\tau} = R, r = yR, \) and, \( x_{\nu_e}^{\text{HT}} = x_{\text{HT}} \) is given by Eq. \( \text{3.17} \) For \( M \rightarrow 0 \) and \( J \rightarrow 0 \), the above function becomes \( F(0) = 1 \) corresponding to the Newtonian background. Putting \( J = 0 \) in Eq. \( \text{4.5} \) gives the function \( F \) for the Schwarzschild background with \( x_{\nu_e}^{\text{Sch}} = \left[ 1 - \frac{R^2}{r^2} \right]^{1/2} \). Eq. \( \text{4.4} \) is for a single neutrino flavor and thus one needs to calculate the total energy deposition contributed by all the three flavors. If the Z’ mediated process is included then Eq. \( \text{4.4} \) will be modified. However, Eq. \( \text{4.5} \) remains unchanged.
for all the three spacetimes mentioned above, as it does not depend on the BSM physics. From Eq. 4.4 assuming that the neutrinos are emitted from the neutrinosphere, D=1.23, the neutrino energy at infinity is \( \sim 10^{52} \text{erg} \), and \( R = 20 \text{ km} \), we obtain the maximum energy deposited as

\[
Q_{\text{Sch}} = 2.4 \times 10^{48} F\left(\frac{M}{R}\right) R^{-3/2} \text{erg.} \tag{4.6}
\]

In the Newtonian background, the SM contribution to \( \nu \bar{\nu} \rightarrow e^+e^- \) process gives the value of the maximum energy deposited as \( Q_{\text{SM}}^{\text{Newt}} = 8.48 \times 10^{47} \text{erg} \). Similarly, for the Schwarzschild background we obtain

\[
Q_{\text{SM}}^{\text{Sch}} \sim 2.5 \times 10^{49} \text{erg}, \tag{4.7}
\]

for \( F\left(\frac{M}{R}\right) \sim 30 \) and \( \frac{R}{M} = 3 \). Hence, the observed GRB energy cannot be explained by modifying the background spacetime with Schwarzschild geometry. If we include rotation in the background spacetime, the rate of energy deposition decreases as compared to the Schwarzschild case.

V. CONSTRAINTS ON \( Z' \) FROM NEUTRINO HEATING IN GRBS: RESULT AND ANALYSIS

In this section, we study the effect of \( Z' \) mediated contribution to the pair annihilation process and calculate the energy deposited for three different spacetimes. We assume that the maximum energy of a GRB is \( \sim O(10^{52} \text{erg}) \), and this whole amount comes from the neutrino pair annihilation process including the contribution from SM (W and Z) and \( Z' \). We also assume that the efficiency of the energy transfer to the GRB is 100%.

A. Constraints on \( Z' \) parameters in Newtonian spacetime from GRB

In this section, we consider the background spacetime for merging neutron stars to be Newtonian. In such a scenario, the SM contribution to \( \nu \bar{\nu} \rightarrow e^+e^- \) process gives the value of the energy deposited as \( \sim 8.48 \times 10^{47} \text{erg} \) whereas the energy required in a GRB is \( 10^{52} \text{erg} \). We investigate if inclusion of \( Z' \) mediated processes can account for the energy needed to power the GRB. In Fig. 2 the pink line shows the variation of the ratio of rate of energy
FIG. 2: Variation of the ratio of rate of energy depositions in BSM and SM processes with respect to \( \frac{M_{Z'}}{g'} \) in Newtonian background. The red line denotes the enhancement factor in energy needed to explain the observed GRB from merging neutron stars.

depositions in BSM to that in SM with respect to \( \frac{M_{Z'}}{g'} \). The figure shows that the ratio decreases with increasing \( \frac{M_{Z'}}{g'} \). As \( \frac{M_{Z'}}{g'} \rightarrow \infty \) (very large), the BSM effect goes away as expected and the ratio becomes unity (here it happens at \( \frac{M_{Z'}}{g'} \sim 10^3 \)). The red line denotes the enhancement factor in energy needed to explain the observed energy in a GRB from merging neutron stars. The point where the pink line intersects the red line is the value of \( \frac{M_{Z'}}{g'} (= 32.59 \text{ GeV}) \) for which the observed GRB energy can be obtained. We note that, for lower values of \( \frac{M_{Z'}}{g'} \) an enhancement more than \( \sim 10^4 \) is possible. Hence, even if the efficiency of the neutrino heating process is less than 100\%, it is still possible to get some parameter space to reach the observed GRB energy.

In Fig. 3, the unshaded region shows the values of the parameters in the \( g' \) vs. \( M_{Z'} \) plane for which the required energy for a GRB can be accessed from neutrino heating in a Newtonian background. The red shaded region is excluded since the observed GRB energy cannot be explained from this combinations of \( g' \) and \( M_{Z'} \).
FIG. 3: Excluded region from neutrino heating in GRB from merging neutron stars in Newtonian background.

B. Constraints on $Z'$ parameters in Schwarzschild spacetime from GRB

As discussed earlier, including the GR effects can enhance the rate of energy deposition via neutrino annihilation. This gives an enhancement factor $\sim 30$ [6]. In this section, we explore the effect of BSM physics if the background spacetime is Schwarzschild. In Fig. 4 we have shown the variation of the ratio of energy deposition in BSM+Schwarzschild to SM+Newtonian cases with respect to $\frac{M_{Z'}}{g'}$. We have shown the variations for $\frac{R}{M} = 3$ (pink) and $\frac{R}{M} = 5$ (brown) which are typical values for merging neutron stars. From the figure it is seen that the ratio decreases with increasing values of $\frac{R}{M}$ and $\frac{M_{Z'}}{g'}$. The red line denotes the enhancement factor in energy deposition required over the Newtonian+SM case to explain the observed GRB from merging neutron stars. The points where pink and brown lines meet the red line correspond to the $\frac{M_{Z'}}{g'}$ values where the observed GRB energy can be attained. The corresponding values are $\frac{M_{Z'}}{g'} = 76.65$ GeV for $\frac{R}{M} = 3$ and $\frac{M_{Z'}}{g'} = 47.98$ GeV for $\frac{R}{M} = 5$. For a very large value of $\frac{M_{Z'}}{g'}$ the ratio becomes constant. Note that, in this limit the BSM effects are negligible and enhancement obtained is $\sim 4 - 30$ (depending on the values of $\frac{R}{M}$).
FIG. 4: Variation of the ratio of rate of energy depositions in BSM+Schwarzschild and SM+Newtonian processes with respect to $\frac{M_{Z'}}{g'}$. We have shown the variations for $\frac{R}{M} = 3$ (pink) and $\frac{R}{M} = 5$ (brown). The red line denotes the enhancement factor in energy deposition required over the Newtonian+SM case to explain the observed GRB from merging neutron stars.

because of changing the spacetime geometry from Newtonian to Schwarzschild.

In Fig. 5 we have shown the variation of the ratio of rate of energy depositions in BSM+Schwarzschild and SM+Newtonian cases with respect to $\frac{R}{M}$. We have varied $\frac{R}{M}$ from 3 to 5 which is relevant for the merging neutron stars. We have shown the variations for $\frac{M_{Z'}}{g'} = 40 \text{ GeV}$ (red), $\frac{M_{Z'}}{g'} = 60 \text{ GeV}$ (pink) and $\frac{M_{Z'}}{g'} \to \infty$ (blue). The ratio decreases as we increase the values of $\frac{M_{Z'}}{g'}$ and $\frac{R}{M}$. This is expected since larger values of $M_{Z'}/g'$ tends towards the SM case and larger $\frac{R}{M}$ values correspond to the Newtonian case. The extra contribution required for the Newtonian+SM scenario to explain the GRB energy is shown by the brown line. The point at which the Schwarzschild+BSM case can explain the GRB energy corresponds to $\frac{M_{Z'}}{g'} = 60 \text{ GeV}$ and $\frac{R}{M} = 3.6$. If there is no BSM physics included, the Schwarzschild background enhances the energy deposition by a factor $\sim 30$ over the Newtonian background (blue line) at $\frac{R}{M} = 3$. 
FIG. 5: Variation of the ratio of rate of energy depositions in BSM+Schwarzschild and SM+Newtonian processes with respect to $\frac{R}{M}$. We have shown the variations for $\frac{M_Z'}{g'} = 40$ GeV (red), $\frac{M_Z'}{g'} = 60$ GeV (pink) and $\frac{M_Z'}{g'} \to \infty$ (blue). The brown line denotes the extra contribution required for the Newtonian+SM scenario to explain the observed GRB energy from merging neutron stars.

In Fig. 6 the shaded regions show the parameter values in the $g'$ vs $M_Z'$ plane for which the required energy for a GRB cannot be accessed from neutrino heating in a Schwarzschild background even after inclusion of $Z'$ mediated diagrams. This is shown for two values of $\frac{R}{M}$. It is shown that smaller values of $\frac{R}{M}$ gives stronger bounds on $g'$.

C. Constraints on $Z'$ parameters in Hartle-Thorne spacetime from GRB

The Hartle-Thorne metric which includes the rotation can also enhance the energy deposition in neutrino heating as compared to the Newtonian background. However, the enhancement due to the inclusion of rotation is smaller as compared to the enhancement due to GR (Schwarzschild) correction. In Fig. 7 we have shown the variation of the ratio of energy depositions in merging neutron stars in Hartle-Thorne and Newtonian background
FIG. 6: Excluded region from neutrino heating in GRB from merging neutron stars in Schwarzschild background.

FIG. 7: (a) Variation of the ratio of energy depositions in merging neutron stars in Hartle-Thorne and Newtonian background with respect to $\frac{M_Z' g'}{g'}$ for fixed $\frac{R}{M}$. (b) Variation of the ratio of energy depositions in merging neutron stars in Hartle-Thorne and Newtonian background with respect to $\frac{M_Z' g'}{g'}$ for fixed $\frac{J}{M^2}$.

with respect to $\frac{M_Z' g'}{g'}$ for fixed $\frac{R}{M} = 3$ (Fig. 7(a)) and fixed $\frac{J}{M^2} = 0.1$ (Fig. 7(b)). In both
the cases, the energy deposition ratio decreases with increasing $\frac{R}{M}$ and $\frac{J_{M^2}}{M^2}$. The red line denotes the enhancement in energy required with respect to the Newtonian+SM scenario to explain the observed GRB from merging neutron stars. In Fig. 7(a), the pink and brown lines denote the variation of the energy deposition ratio for the values $\frac{J_{M^2}}{M^2} = 0.1$ and 0.8 respectively for fixed $\frac{R}{M} = 3$. Similarly, in Fig. 7(b) the pink and brown lines denote the variation of the energy deposition ratio for the values $\frac{R}{M} = 3$ and 5 respectively for fixed $\frac{J_{M^2}}{M^2} = 0.1$. In Fig. 7(a) the points where the red line intersects the pink and brown lines correspond to the values of $\frac{M_{Z'_{g'}}}{g'} = 72.57$ GeV and 61.49 GeV, and $\frac{J_{M^2}}{M^2} = 0.1, 0.8$ respectively for $\frac{R}{M} = 3$ where the observed GRB energy can be explained in HT spacetime including $Z'$ mediated processes. Similarly, in Fig. 7(b) the points where the red line intersects the pink and brown lines correspond to the values of $\frac{M_{Z'_{g'}}}{g'} = 72.57$ GeV and 46.69 GeV, and $\frac{R}{M} = 3, 5$ respectively for $\frac{J_{M^2}}{M^2} = 0.1$ where the observed GRB energy can be explained in HT spacetime. One can lower the value of $\frac{M_{Z'_{g'}}}{g'}$ with increasing $\frac{R}{M}$ and $\frac{J_{M^2}}{M^2}$ to explain the observed GRB energy. In Fig. 8 we have shown the variation of the ratio of energy depositions in

![Graph](image-url)

(a) $\dot{Q}_{\text{BSM}}^\text{HT}$ vs. $\frac{R}{M}$, $\frac{J_{M^2}}{M^2} = 0.1$

(b) $\dot{Q}_{\text{BSM}}^\text{HT}$ vs. $\frac{R}{M}$, $\frac{J_{M^2}}{M^2} = 0.8$

FIG. 8: (a) Variation of the ratio of energy depositions in merging neutron stars in Hartle-Thorne and Newtonian background with respect to $\frac{R}{M}$ for $\frac{J_{M^2}}{M^2} = 0.1$. (b) Variation of the ratio of energy depositions in merging neutron stars in Hartle-Thorne and Newtonian background with respect to $\frac{R}{M}$ for $\frac{J_{M^2}}{M^2} = 0.8$. 

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merging neutron stars in Schwarzschild and Newtonian background with respect to $\frac{R}{M}$ for two fixed values of $\frac{J}{M^2} = 0.1, 0.8$ and with different values of $\frac{M_{Z'}}{g'}$. The red, pink, and blue lines denote the variations with $\frac{R}{M}$ for $\frac{M_{Z'}}{g'} = 40$ GeV, 60 GeV, $\infty$ respectively. The energy deposition ratio decreases with increasing $\frac{M_{Z'}}{g'}$. The brown line denotes the excess amount of energy required to explain the observed GRB with respect to the Newtonian+SM picture. The points where the brown lines intersect the pink lines in both the plots correspond to the values $\frac{R}{M} = 3.5, \frac{J}{M^2} = 0.1$ (Fig. 8(a)) and $\frac{R}{M} = 3.01, \frac{J}{M^2} = 0.8$ (Fig. 8(b)) respectively where the observed GRB energy is achieved in HT+BSM case.

In Fig. 9 we have shown the shaded region in $g'$ vs $M_{Z'}$ plane for which the required heating is not achieved in the Hartle-Thorne background even after including the effect of $Z'$ mediated process for different values of $\frac{R}{M}$ and $\frac{J}{M^2}$. Lower values of $\frac{R}{M}$ and $\frac{J}{M^2}$ gives stronger bounds on $g'$. 

FIG. 9: Excluded region from neutrino heating in GRB from merging neutron stars in Hartle-Thorne background.
FIG. 10: The allowed parameter space in $g' - M_{Z'}$ plane. The existing constraints are from BaBar (brown), LHCb (green), Beam-Dump (purple), and $\nu - e$ scattering (cyan). The region above the blue line denotes the exclusion region from GRB from neutrino pair annihilation in the Newtonian background. Similarly the region above the black line denotes the exclusion region in the Schwarzschild background. The regions above the magenta and red lines denote the exclusion regions in Hartle-Thorne background with low and high $J M^2$ respectively.

VI. COMBINED ANALYSIS IN CONSTRAINING $g'$ AND $M_{Z'}$

In Fig. 10 we superimpose the existing constraints from BaBar [56] (brown), LHCb [57] (green), Beam-Dump (purple) [58, 61], and $\nu - e$ scattering [38, 39] (cyan) on those obtained from neutrino heating of a GRB for different spacetime backgrounds. The region above the blue line denotes the exclusion region from GRB from neutrino pair annihilation in the Newtonian background. Similarly the region above the black line denotes the exclusion region in the Schwarzschild background. The regions above the magenta and red lines denote the exclusion regions in Hartle-Thorne background with low $J M^2 = 0.1$ and high $J M^2 = 0.8$ respectively. The bound on the gauge coupling in Schwarzschild background is stronger than that for Newtonian and Hartle-Thorne background. The bounds on the gauge coupling get
stronger with decreasing the values of $\frac{R}{M}$ and $\frac{J}{M^2}$. The constraints on the gauge coupling from GRB is the weakest in the Newtonian background. Consideration of the neutrino heating of GRB due to $Z'$ mediated processes can disfavor low $M_Z'$-high $g'$ values, which cannot be accessed for instance by the BaBar and LHCb experiments. However, this parameter space can be probed from neutrino-electron scattering experiments. Combining all these constraints, the unshaded region with $M_{Z'} \sim (0.01 - 10) \text{ GeV}$ and $g' \sim (10^{-6} - 10^{-4})$ remains allowed for which the $Z'$ mediated process can successfully enhance the efficiency of neutrino heating to power a GRB.

VII. CONCLUSION

The neutrino annihilation process ($\nu\bar{\nu} \rightarrow e^+e^-$) is important in stellar environments since it can deposit energy which can aid type II supernovae explosions, merging neutron stars, GRBs etc. The inclusion of general relativistic corrections by changing the Newtonian spacetime to Schwarzschild can enhance the rate of energy deposition. Inclusion of rotations i.e; introducing the Hartle-Thorne metric can also increase the energy deposition as compared to the Newtonian case. However, in all the above cases, the energy deposited is not sufficient to power a GRB ($\sim 10^{52} \text{erg}$). In this paper, we investigate the scope of beyond standard model contribution to the neutrino pair annihilation process in achieving the observed energy required for the GRB. In particular, we concentrate on the scenario where the standard model is augmented by $U(1)_{B-L}$ gauge group. This scenario gives rise to additional $Z'$ mediated diagrams. We obtain the values on the mass ($M_{Z'}$) of the extra $Z'$ and coupling ($g'$) for which addition of such processes can contribute enough energy in powering GRBs. We also show the excluded regions in the $g' - M_{Z'}$ plane for which such enhancement is not viable. We present our results for Newtonian, Schwarzschild, and Hartle-Thorne backgrounds and study the combined effect of modifying the background spacetime and inclusion of BSM physics. From the GRB data, we obtain the parameter values for which the successful energy deposition is possible, as $g' \lesssim 2.98 \times 10^{-4}$ for Newtonian background, $g' \lesssim 1.28 \times 10^{-4}$ for $\frac{R}{M} = 3$ in Schwarzschild background, $g' \lesssim 1.35 \times 10^{-4}$ for $\frac{J}{M^2} = 0.1$, $\frac{R}{M} = 3$, and $g' \lesssim 1.62 \times 10^{-4}$ for $\frac{J}{M^2} = 0.8$, $\frac{R}{M} = 3$ in Hartle-Thorne background for the gauge boson with mass in the
range $0.01 \text{ GeV} \lesssim M_{Z'} \lesssim 10 \text{ GeV}$. The stronger bound on the gauge coupling comes for the Schwarzschild case. We compare the bounds with those obtained in BaBar, LHCb, and $\nu$-electron scattering experiments. Our results show that imposing the constraints on GRB heating can access lower values of $M_{Z'}$ as compared to BaBar and LHCb. However, this regions can also be probed from neutrino-electron scattering experiments like TEXONO, GEMMA, and Borexino. It can be seen that some of the regions where the energy required for a GRB can be obtained is disfavoured by the laboratory experiments. However, a small area remains where the $Z'$ mediated process can successfully enhance the energy deposition due to neutrino annihilation such that it can contribute to the GRB energy. This is given as $M_{Z'} \sim (0.01 - 10) \text{ GeV}$ and $g' \sim (10^{-6} - 10^{-4})$. In conclusion, our study highlights the importance of inclusion of BSM contributions to the neutrino annihilation processes in energizing GRBs. This can open up new avenues in understanding the physics of GRBs.

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