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The final outcome of dissipative collapse in the presence of Λ

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Abstract. We investigate the role played by the cosmological constant during gravitational collapse of a radiating star with vanishing Weyl stresses in the interior. We highlight the role played by the cosmological constant during the latter stages of collapse. The evolution of the temperature of the collapsing body is studied by employing causal heat transport equation. We show that the inclusion of the cosmological constant enhances the temperature within the stellar core.

Keywords. Gravitational collapse; black holes; cosmological constant.

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1. Introduction

The end state of a collapsing star remains an important problem in theoretical astrophysics. The Cosmic Censorship Conjecture due to Penrose held sway until the discovery of solutions admitting naked singularities. Relativistic stellar models incorporating physically viable matter distributions such as heat flux, radiation, anisotropic stresses and the electromagnetic field have made the study of gravitational collapse more tractable. The simple dust ball collapse first studied by Oppenheimer and Snyder [1] has evolved into a more general treatment of the collapsing sphere in which gravitational and thermodynamical effects determine the outcome of collapse. The Vaidya solution [2] has made it possible to study radiating stars with a nonempty exterior. The junction conditions for a radiating spherically symmetric star was presented by Santos [3]. These conditions have led to a rich class of radiating stars which continue to provide insight into the problem of dissipative collapse. General treatments of the influence of thermodynamical fluxes during collapse as well as physically motivated stellar models such as Euclidean stars and expansion-free collapse have been provided by Herrera and co-workers [4,5]. Their

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ground-breaking work has laid the foundations for a more realistic approach to the study of gravitational collapse of radiating stars.

The cosmological constant problem has presented a strong challenge to both observational as well as theoretical physics. Present day observations from Type Ia supernovae data, baryonic acoustic oscillations and high redshift data point to a small, positive cosmological constant. Quantum field theory, on the other hand, predicts a large theoretical value for the cosmological constant. This constant has been associated with dark energy or gravitational aether. The uncertainty in the equation of state at nuclear densities (which may be the case at various epochs of a star's evolution) begs the question as to what degree would a nonzero cosmological constant affect the evolution of a collapsing stellar configuration. In this study we investigate the role played by the cosmological constant in the gravitational collapse of conformally flat, radiating spheres. The influence of the cosmological constant on bounded matter distributions has been studied in various scenarios. Rindler and Ishak [6] showed that a positive cosmological constant diminishes the classical bending of light by a localized, spherically symmetric mass distribution. Comprehensive studies of static fluid spheres in the presence of a cosmological constant have led to a modification of Buchdahl's compactness ratio $M/R \leq 4/9$ (see Andreasson and Boehmer [7] and references therein). Recently, Chan *et al* [8] investigated the influence of the cosmological constant on gravastar formation. Their model consisted of a de Sitter interior space-time, matched to an infinitely thin fluid shell with a barotropic equation of state, which was in turn matched to an external de Sitter–Schwarzschild space-time. They showed that the formation of these particular models is affected by the relative magnitudes of the interior and exterior cosmological constants. In their study of the gravitational collapse of null strange quark fluid and its influence on cosmic censorship, Ghosh and Dadhich [9] indicate that the bag constant B appearing in the equation of state $p = \frac{1}{n}(\rho - 4B)$ makes a similar contribution as the cosmological constant to the dynamics of the collapsing fluid. Physically, the inclusion of the strange quark matter component favours the formation of black holes. Govender and Thirukkanesh [10] provided a class of radiating stellar models with heat dissipation in the presence of a cosmological constant. A study of the temperature profiles of these models indicated that the cosmological constant enhances the temperature at each interior point within the stellar core.

The general solution for the interior of a spherically symmetric, conformally flat radiating star was first provided by Banerjee *et al* [11]. A simple radiating model with heat flux was presented in which the exterior space-time was described by the outgoing Vaidya solution. In a more recent study by Herrera *et al* [4], the general conformally flat model was resurrected in which they were able to solve the boundary condition required for the smooth matching of the interior and exterior space-times for particular cases. The physical viability of these models was studied by Maharaj and Govender [12]. Subsequently, Herrera *et al* [5] and Misthry *et al* [13] obtained further classes of conformally flat radiating stars. In this paper we generalize the results of Misthry *et al* to include the cosmological constant. We further provide an analysis of the physical behaviour of our model within the framework of extended irreversible thermodynamics. The influence of the cosmological constant on the temperature profiles is clearly exhibited in both the causal and noncausal theories. We draw comparisons with the models of Misthry *et al* and confirm earlier findings by Govender and Thirukkanesh.

2. Field equations

We investigate a spherically symmetric radiating star with shear-free matter. The interior metric for shear-free matter is given by

$$ds^2 = -A^2 dt^2 + B^2[dr^2 + r^2(d\theta^2 + \sin^2\theta d\phi^2)], \quad (1)$$

where A and B are functions of both the temporal coordinate t and the radial coordinate r. The energy–momentum tensor for the interior matter distribution is described by

$$T_{ab} = (\rho + p)u_a u_b + p g_{ab} + q_a u_b + q_b u_a. \quad (2)$$

For the line element (1) and matter distribution (2), the coupled Einstein field equations

$$R_{ab} - \frac{1}{2} R g_{ab} + \Lambda g_{ab} = T_{ab} \quad (3)$$

become

$$\rho = 3\frac{1}{A^2}\frac{\dot{B}^2}{B^2} - \frac{1}{B^2}\left(2\frac{B''}{B} - \frac{B'^2}{B^2} + \frac{4}{r}\frac{B'}{B}\right) - \Lambda, \quad (4)$$

$$p = \frac{1}{A^2}\left(-2\frac{\ddot{B}}{B} - \frac{\dot{B}^2}{B^2} + 2\frac{\dot{A}}{A}\frac{\dot{B}}{B}\right)$$
$$+ \frac{1}{B^2}\left(\frac{B'^2}{B^2} + 2\frac{A'}{A}\frac{B'}{B} + \frac{2}{r}\frac{A'}{A} + \frac{2}{r}\frac{B'}{B}\right) + \Lambda, \quad (5)$$

$$p = -2\frac{1}{A^2}\frac{\ddot{B}}{B} + 2\frac{\dot{A}}{A^3}\frac{\dot{B}}{B} - \frac{1}{A^2}\frac{\dot{B}^2}{B^2} + \frac{1}{r}\frac{A'}{A}\frac{1}{B^2}$$
$$+ \frac{1}{r}\frac{B'}{B^3} + \frac{A''}{A}\frac{1}{B^2} - \frac{B'^2}{B^4} + \frac{B''}{B^3} + \Lambda, \quad (6)$$

$$q = -\frac{2}{AB^2}\left(-\frac{\dot{B}'}{B} + \frac{B'\dot{B}}{B^2} + \frac{A'}{A}\frac{\dot{B}}{B}\right), \quad (7)$$

where the heat flux $q^a = (0, q, 0, 0)$ has only the nonvanishing radial component. The system of equations (4)–(7) governs the general situation in describing matter distributions with isotropic pressures in the presence of heat flux and cosmological constant Λ for a spherically symmetric relativistic stellar object. From (4)–(7), we observe that if the gravitational potentials $A(t, r)$ and $B(t, r)$ are known, then the expressions for the matter variables ρ, p and q follow immediately. The system (4)–(7) contains four equations with five unknowns A, B, p, q and ρ so that we may specify one of the variables to solve the system. The components of the Weyl tensor for the line element (1) are

$$C_{2323} = -r^4\left(\frac{B}{A}\right)^2 \sin^2\theta C_{0101}, \quad (8)$$
$$= 2r^2\left(\frac{B}{A}\right)^2 \sin^2\theta C_{0202}, \quad (9)$$
$$= 2r^2\left(\frac{B}{A}\right)^2 C_{0303}, \quad (10)$$
$$= -2r^2 \sin^2\theta C_{1212}, \quad (11)$$
$$= -2r^2 C_{1313}, \quad (12)$$

where

$$C_{2323} = \frac{r^4}{3} B^2 \sin^2\theta \left[\left(\frac{A'}{A} - \frac{B'}{B} \right) \left(\frac{1}{r} + 2\frac{B'}{B} \right) - \left(\frac{A''}{A} - \frac{B''}{B} \right) \right]. \tag{13}$$

In order to ensure the vanishing of the Weyl stresses in the interior space-time we must have $C_{2323} = 0$ which yields

$$A = (C_1(t) r^2 + 1) B. \tag{14}$$

Equating (5) and (6), and using (14) we obtain the condition of pressure isotropy

$$\frac{B''}{B'} - 2\frac{B'}{B} - \frac{1}{r} = 0. \tag{15}$$

On integrating (15) we get

$$B = \frac{1}{C_2(t) r^2 + C_3(t)}. \tag{16}$$

3. Junction conditions

The exterior space-time of our model is described by the Vaidya solution with nonvanishing cosmological constant Λ

$$\mathrm{d}s^2 = -\left(1 - \frac{2m(v)}{r} - \frac{1}{3}\Lambda r^2\right) \mathrm{d}v^2 - 2\mathrm{d}v\mathrm{d}r + r^2 \left(\mathrm{d}\theta^2 + \sin^2\theta \mathrm{d}\phi^2\right), \tag{17}$$

where $m(v)$ denotes the mass of the fluid as measured by an observer at infinity. To obtain a complete description of a radiating star, the interior space-time must be smoothly matched to the exterior space-time. The junction conditions have recently been presented by Govender and Thirukkanesh [10] and we only present the main results here.

$$A(r_\Sigma, t)\mathrm{d}t = \left(1 - \frac{2m}{r_\Sigma} - \frac{1}{3}\Lambda r^2{}_\Sigma + 2\frac{\mathrm{d}r_\Sigma}{\mathrm{d}v}\right)^{1/2} \mathrm{d}v \tag{18}$$

$$r_\Sigma B(r_\Sigma, t) = r_\Sigma(v) \tag{19}$$

$$m(v) = \left(\frac{r^3 B}{2A^2} \dot{B}^2 - r^2 B' - \frac{r^3}{2B} B'^2 - \frac{1}{6}\Lambda r^3 B^3 \right)_\Sigma \tag{20}$$

$$p_\Sigma = (qB)_\Sigma. \tag{21}$$

To fix the temporal behaviour of our model, junction condition (21) needs to be solved. For our line element (1) and the assumption of vanishing Weyl stresses, eq. (21) reduces to the nonlinear equation

$$\begin{aligned} &\ddot{C}_2 b^2 + \ddot{C}_3 - \frac{3}{2}\frac{(\dot{C}_2 b^2 + \dot{C}_3)^2}{C_2 b^2 + C_3} - \frac{\dot{C}_1 b^2 (\dot{C}_2 b^2 + \dot{C}_3)}{C_1 b^2 + 1} - 2(C_1 \dot{C}_3 - \dot{C}_2) b \\ &\quad + 2\frac{(C_1 b^2 + 1)}{C_2 b^2 + C_3} [C_2 (C_2 - 2C_1 C_3) b^2 + C_3 (C_1 C_3 - 2C_2)] \\ &\quad + \frac{\Lambda}{2}\frac{(C_1 b^2 + 1)^2}{C_2 b^2 + C_3} = 0, \end{aligned} \tag{22}$$

where $r = b$ determines the boundary of the star. Following the ansatz adopted by Misthry *et al* [13], eq. (22) can be recast into a simpler form by introducing the transformation

$$U = C_1 b^2 + 1, \tag{23}$$

where $U = U(t)$. Using this transformation eq. (22) takes the form

$$\begin{aligned}(\dot{C}_2 b^2 + \dot{C}_3)\dot{U} &+ \left[\frac{3}{2}\frac{(\dot{C}_2 b^2 + \dot{C}_3)^2}{C_2 b^2 + C_3} - \frac{2}{b}(\dot{C}_2 b^2 + \dot{C}_3) - (\ddot{C}_2 b^2 + \ddot{C}_3)\right] U \\ &+ 2\left[\frac{\dot{C}_3}{b} - \frac{1}{C_2 b^2 + C_3}\left(C_2^2 b^2 - \frac{C_3^2}{b^2}\right)\right] U^2 \\ &+ \left[\left(2C_2 b^2 - C_3\right)\frac{C_3}{b^2} - \frac{\Lambda}{4}\right]\frac{2U^3}{C_2 b^2 + C_3} = 0. \end{aligned} \tag{24}$$

Note that eq. (24) is an Abel's equation of the first kind in U. This can we written as

$$\mathscr{A}\dot{U} + \mathscr{B}U + \mathscr{C}U^2 + \mathscr{D}U^3 = 0, \tag{25}$$

where

$$\begin{aligned}\mathscr{A} &= \dot{C}_2 b^2 + \dot{C}_3, \\ \mathscr{B} &= \frac{3}{2}\frac{(\dot{C}_2 b^2 + \dot{C}_3)^2}{C_2 b^2 + C_3} - \frac{2}{b}(\dot{C}_2 b^2 + \dot{C}_3) - (\ddot{C}_2 b^2 + \ddot{C}_3), \\ \mathscr{C} &= 2\left[\frac{\dot{C}_3}{b} - \frac{1}{C_2 b^2 + C_3}\left(C_2^2 b^2 - \frac{C_3^2}{b^2}\right)\right], \\ \mathscr{D} &= \left[\left(2C_2 b^2 - C_3\right)\frac{C_3}{b^2} - \frac{\Lambda}{4}\right]\frac{2}{C_2 b^2 + C_3}. \end{aligned}$$

In general, Abel's equations are difficult to solve. However, we obtain classes of solutions by imposing restrictions on $\mathscr{A}$, $\mathscr{B}$, $\mathscr{C}$ and $\mathscr{D}$ in the following sections. For the case $\mathscr{A} = 0$ we obtain

$$C_2 b^2 + C_3 = \alpha,$$

where α is an arbitrary constant. Therefore eq. (24) becomes an algebraic equation

$$\left[\frac{\dot{C}_3}{b} - \frac{1}{\alpha}\left(C_2^2 b^2 - \frac{C_3^2}{b^2}\right)\right] U^2 + \left[\left(2C_2 b^2 - C_3\right)\frac{C_3}{b^2} - \frac{\Lambda}{4}\right]\frac{U^3}{\alpha} = 0. \tag{26}$$

This then yields the solution

$$C_1 = \begin{cases} \dfrac{-1}{b^2}, & U = 0; \\ \dfrac{4\alpha}{4C_3(2\alpha - 3C_3) - \Lambda b^2}\left[\dfrac{\alpha}{b^2} - \dfrac{4C_3}{b^2} + \dfrac{3C_3^2}{\alpha b^2} - \dfrac{\dot{C}_3}{b} + \dfrac{\Lambda}{4\alpha}\right], & U \neq 0 \end{cases} \tag{27}$$

$$C_2 = \frac{\alpha - C_3}{b^2} \tag{28}$$

$$C_3 = \text{arbitrary function of } t. \tag{29}$$

Note that we obtain this class of solutions in terms of arbitrary function C_3 without any integration. When $\Lambda = 0$ the class of solutions (27)–(29) reduces to the first category of solutions found by Misthry *et al* [13]. If we set $\mathscr{C} = 0$ we get

$$\frac{\dot{C}_3}{b} - \frac{1}{C_2b^2 + C_3}\left(C_2^2b^2 - \frac{C_3^2}{b^2}\right) = 0.$$

This is a quadratic equation in C_2 and we have

$$C_2 = \frac{\dot{C}_3 b \pm \sqrt{\dot{C_3}^2 b^2 + 4C_3(b\dot{C}_3 + C_3)}}{2b^2}.$$

Hence if we specify C_3, we obtain C_2. In this case the Abel equation (24) becomes

$$\begin{aligned}(\dot{C}_2b^2 + \dot{C}_3)\dot{U} \ &+ \left[\frac{3}{2}\frac{(\dot{C}_2b^2 + \dot{C}_3)^2}{C_2b^2 + C_3} - \frac{2}{b}(\dot{C}_2b^2 + \dot{C}_3) - (\ddot{C}_2b^2 + \ddot{C}_3)\right]U \\ &= -\left[\left(2C_2b^2 - C_3\right)\frac{C_3}{b^2} - \frac{\Lambda}{4}\right]\frac{2}{C_2b^2 + C_3}U^3. \qquad (30)\end{aligned}$$

On integrating (30), we obtain

$$U = \frac{e^{2t/b}(\dot{C}_2b^2 + \dot{C}_3)}{(C_2b^2 + C_3)^{3/2}\left[K_2 + \frac{1}{b^2}\int \frac{e^{4t/b}[4C_3(2C_2b^2 - C_3) - \Lambda b^2](\dot{C}_2b^2 + \dot{C}_3)}{(C_2b^2 + C_3)^4}\mathrm{d}t\right]^{1/2}}, \qquad (31)$$

where K_2 is an integrating constant. Therefore, we have the solution

$$C_1 = \frac{1}{b^2}\left[\frac{e^{2t/b}(\dot{C}_2b^2 + \dot{C}_3)}{(C_2b^2+C_3)^{3/2}\left[K_2 + \frac{1}{b^2}\int \frac{e^{4t/b}[4C_3(2C_2b^2 - C_3) - \Lambda b^2](\dot{C}_2b^2 + \dot{C}_3)}{(C_2b^2 + C_3)^4}\mathrm{d}t\right]^{1/2}} - 1\right] \qquad (32)$$

$$C_2 = \frac{\dot{C}_3 b \pm \sqrt{\dot{C_3}^2 b^2 + 4C_3(\dot{C}_3 b + C_3)}}{2b^2} \qquad (33)$$

$$C_3 = \text{arbitrary function of } t. \qquad (34)$$

Again infinite classes of solutions are possible for suitable choices of C_3. When the cosmological constant Λ vanishes, the solutions (32)–(34) reduce to the third category of solutions presented by Misthry *et al*. We have successfully demonstrated the existence of an infinite class of solutions describing a radiating, collapsing sphere with vanishing Weyl stresses within its interior in the presence of cosmological constant. The cases $\mathscr{B} = 0$ and $\mathscr{D} = 0$ follow in a similar manner as worked out by Misthry *et al*.

4. Possible end-states

We are now in a position to investigate the role played by the cosmological constant in the final outcome of dissipative collapse in a particular radiating model. To this end we will consider the solution given by (27)–(29) for the case $U \neq 0$. The mass function (20) yields

$$m = \left(\frac{r}{6^{1/3}(\frac{r^2(\alpha - C_3)}{b^2} + C_3)}\right)^3 \left[-\Lambda + \frac{12\alpha C_3}{b^2} - \frac{12C_3^2}{b^2}\right]$$
$$+\left(\frac{r}{6^{1/3}(\frac{r^2(\alpha - C_3)}{b^2} + C_3)}\right)^3$$
$$\times\left[\frac{3(b^2-r^2)^2(b^2\Lambda-8\alpha C_3+12C_3^2)^2\dot{C}_3^2}{(-b^4\Lambda+r^2(4\alpha^2+b^2\Lambda)+8(b^2-2r^2)\alpha C_3+12(r^2-b^2)C_3^2-4br^2\alpha\dot{C}_3)^2}\right]. \tag{35}$$

We can now easily calculate the ratio $2M/rB$ evaluated at the boundary which gives us

$$\left(\frac{2M}{rB}\right)_\Sigma = \frac{12(\alpha - C_3)C_3 - b^2\Lambda}{3\alpha^2}, \tag{36}$$

where we have taken the boundary to be $r = b$. It is clear that the time of formation of the horizon is reduced in the presence of a positive cosmological constant. The total luminosity of the collapsing star as perceived by an observer at infinity is given by

$$L_\infty = -\frac{\mathrm{d}M}{\mathrm{d}v} = \frac{b(\alpha - 2C_3)^2(b^2\Lambda - 8\alpha C_3 + 12C_3^2)\dot{C}_3}{2\alpha^4(\alpha - 2C_3 - b\dot{C}_3)} \tag{37}$$

evaluated on Σ. We note that at the time of formation of the black hole, $L_\infty = 0$, i.e., no radiation from the stellar surface reaches our observer at infinity. This occurs when

$$\dot{C}_3 = 0, \qquad C_3 = \frac{\alpha}{2},$$
$$C_3 = \frac{1}{6}\left[2\alpha - \sqrt{4\alpha^2 - 3b^2\Lambda}\right], \qquad C_3 = \frac{1}{6}\left[2\alpha + \sqrt{4\alpha^2 - 3b^2\Lambda}\right]. \tag{38}$$

In the case of vanishing cosmological constant, L_∞ becomes zero when

$$\dot{C}_3 = 0, \qquad C_3 = \frac{\alpha}{2}, \qquad C_3 = 0, \qquad C_3 = \frac{2\alpha}{3}. \tag{39}$$

Note that the case $\dot{C}_3 = 0$ leads to the vanishing of the heat flux. Since the function C_3 represents the temporal evolution of the model, we note that the case $C_3 = \alpha/2$ corresponds to $R_{\mathrm{Sch}} = 2M_\Sigma$ which marks the time when the collapsing sphere crosses its Schwarzschild radius to form a black hole. The case $C_3 = 0$ corresponds to the initial epoch, when the system starts to collapse. Since $C_3(t)$ is an arbitrary function we can choose it such that the evolution of the collapsing system starts off in the remote past at $t = -\infty$ and proceeds towards $t = 0$. A similar analysis for the case of vanishing cosmological constant was performed by Sarwe and Tikekar [14] in which the collapse proceeds from an initial static configuration.

5. Temperature profiles

To exhibit the relaxational effects on the temperature profile, we shall employ the causal transport equation for heat flux given by

$$\tau h_a{}^b \dot{q}_b + q_a = -\kappa (D_a T + T \dot{u}^a), \tag{40}$$

where τ is the relaxation time for the thermal signals. Setting $\tau = 0$ in eq. (40), we regain the so-called Eckart transport equations which predict infinite propagation velocities for the dissipative fluxes. For the line element (1) the causal transport equation (40) reduces to

$$\tau (qB)\dot{} + A(qB) = -\kappa \frac{(AT)'}{B} \tag{41}$$

which governs the behaviour of the temperature. Setting $\tau = 0$ in eq. (41) we obtain the noncausal Fourier heat transport equation

$$A(qB) = -\kappa \frac{(AT)'}{B} \tag{42}$$

which yields reasonable temperature profiles provided the fluid is close to hydrostatic equilibrium.

Employing the thermodynamic coefficients for radiative transfer as motivated by Govender *et al* [15], the thermal conductivity takes the form

$$\kappa = \gamma T^3 \tau_{\mathrm{c}}, \tag{43}$$

where γ (≥ 0) is a constant and τ_{c} is the mean collision time between the massless and massive particles. Following [16] we assume a generalized power-law behaviour for τ_{c}.

$$\tau_{\mathrm{c}} = \left(\frac{\alpha}{\gamma}\right) T^{-\sigma}, \tag{44}$$

where α (≥ 0) and σ (≥ 0) are constants. In our calculations we further assume that the velocity of thermal dissipative signals is comparable to the adiabatic sound speed which is ensured if the relaxation time is proportional to the collision time:

$$\tau = \left(\frac{\beta\gamma}{\alpha}\right) \tau_{\mathrm{c}}, \tag{45}$$

where τ (≥ 0) is a constant. The constant β can be thought of as a causality index, measuring the strength of relaxational effects, with $\beta = 0$ giving the noncausal case. Using the above definitions for τ and κ, eq. (41) takes the form

$$\beta (qB)\dot{} T^{-\sigma} + A(qB) = -\alpha \frac{T^{3-\sigma}(AT)'}{B}. \tag{46}$$

The Eckart temperature is readily obtained by setting $\beta = 0$ in (46). Exact solutions to eq. (46) have been presented by Govinder and Govender [17] for constant collision times as well as variable collision times. For our model we are able to plot the temperature profiles for the special case of constant collision time. This assumption may hold true for a brief period of the collapse process. We make use of solutions (27)–(29)

for the case $U \neq 0$ and $C_3(t) = a\mathrm{e}^t$ where a is a constant. It is clear from figures 1 and 2 that the temperatures are well behaved in both the causal and noncausal theories. More importantly, we note that the temperatures in the presence of the cosmological constant are higher throughout the stellar core compared to the corresponding model of Misthry *et al* [13] ($\Lambda = 0$). This result confirms the earlier findings by Govender and Thirukkanesh [10]. Here our model has nonvanishing acceleration while the model studied by Govender and Thirukkanesh is acceleration-free. It is clear from figure 2 that relaxational effects enhance the temperature at each interior point by a factor of ten. The causal temperature gradient is higher than its noncausal counterpart, confirming the perturbative results presented by Govender *et al* [16]. It would be interesting and relevant to include the effects of shear in our radiating model. Work in this direction has been initiated.

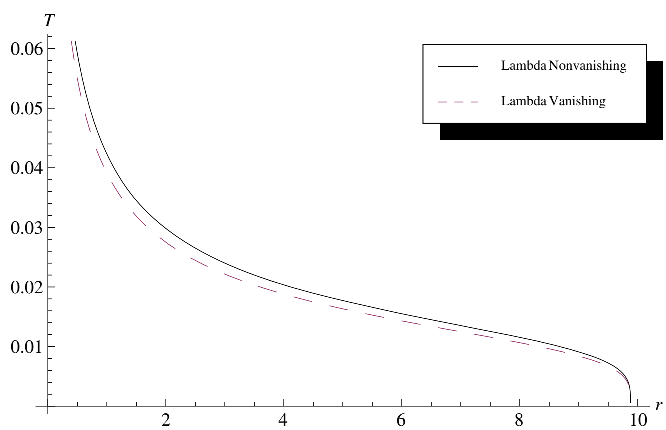

Figure 1. Noncausal temperature profiles.

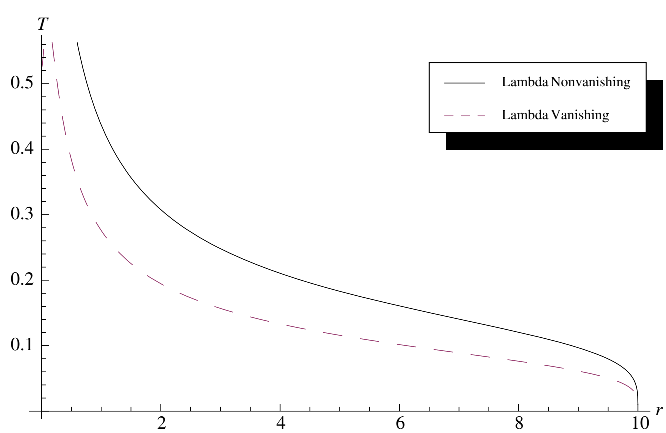

Figure 2. Causal temperature profiles.

6. Conclusion

We have successfully modelled a radiating star undergoing dissipative collapse in the presence of a positive cosmological constant. The temporal evolution equation arising from the junction conditions was solved exactly and we have presented various classes of solutions with nonvanishing cosmological constant. We further studied the end-state of the collapse of a particular model and showed that the time of formation of the horizon is independent of the cosmological constant. The evolution of the temperature of the star was obtained using a causal heat transport equation. We have shown that the presence of the cosmological constant enhances the core temperature. This confirms the results obtained earlier by Govender and Thirukannesh [10].

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