Weak Interactions as a mechanism against Gravitational Collapse in Primordial Baryonic Stars

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Abstract: This work investigates whether weak interactions mediated by Z-boson exchange can provide a stabilising pressure against gravitational collapse (GC) in primordial baryonic stars. Building on previous works, we completed the equation of state by incorporating the degeneracy pressure of an ultra-relativistic fermionic gas. A crossover density $n^* \approx 9 \cdot 10^{10}$ fm^{-3} is identified, marking the transition between weak-interaction-dominated and degeneracy-dominated regimes. The Tolman-Oppenheimer-Volkoff equation is solved numerically with a boundary condition derived from the thermal radiation pressure of the primordial plasma, yielding a universal stellar radius of approximately 15 meters. Stability analysis via the turning point theorem reveals that all stable configurations have central densities well below the crossover density, demonstrating that weak interactions do not play a relevant structural role in these compact objects.
Keywords: Primordial Baryonic Stars, Weak Interactions, Degenerate Fermi gas, Tolman-Oppenheimer-Volkoff Equation.
SDGs: This work is related to SDG 4 and SDG 17.

I. INTRODUCTION

In the primordial universe, fractions of a second after the Big Bang, matter existed at temperatures and densities far exceeding those of any astrophysical object known today. Under such conditions, interactions typically negligible in modern-day stellar physics may have played a non-trivial structural role. In normal baryonic stars, although negligible, weak force interactions act coherently over these objects carrying a net weak charge since, what matters is the total conserved weak charge regardless of its carriers [1]. This coherence, combined with the repulsive nature of the Z-boson exchange between equal weak charges [2, 3], makes it a possible candidate for a stabilising pressure against gravitational collapse at ultra-high densities, such as those found in the early universe.

Espriu [1] showed that solving the Tolman-Oppenheimer-Volkoff equation [4, 5] with the equation of state derived from Z-mediated weak interactions yields stable, compact objects with radii of a few metres and masses up to $\sim 10^{-3}$ solar masses. The equation of state used, of the stiff form $p \approx \varepsilon$ in the ultra-high density limit, is a theoretical extremum proposed by Zeldovich [6], long considered physically unrealisable. The analysis relies on the approximation that the contribution of the degenerate Fermi sea is negligible due to the extreme densities in the interior of the stellar object [7]. This assumption becomes less justifiable in the outer layers of the object, where densities are comparable to those of present-day barionic stars [8].

This thesis aims to extend the framework presented in [1] by incorporating the degeneracy pressure of an ultra-relativistic fermionic gas into the energy density, scaling as $\varepsilon \propto n^{4/3}$ with the number density of weak charge carriers [7, 9], therefore relaxing the equation of state proposed by Zeldovich [6], and to assess whether the stability of these weakly charged compact objects survives this more complete treatment.

II. THEORETICAL BACKGROUND

A. Neutron Stars and Gravitational Collapse

Neutron stars constitute the final stable configuration of large stellar objects. These stars are sustained against gravitational collapse by the degeneracy pressure of a dense fermionic sea of neutrons combined with repulsive nuclear forces [7]. Their maximum mass, known as the Tolman-Oppenheimer-Volkoff limit, is around a few solar masses [4, 5], beyond which the degeneracy pressure and strong interactions are insufficient to resist gravity, leading to a total collapse into a black hole. More precisely, as shown by Hawking and Ellis [10], once a stellar object crosses its Schwarzschild horizon $r_s = 2GM_o/c^2$ a physical singularity at $r = 0$ becomes inevitable within General Relativity.

B. Weak Interactions in Dense Matter

The weak interaction, mediated by Z and $W^{\pm}$ bosons, is characterised by the Fermi constant $G_F \simeq 1.166 \cdot 10^{-5}\text{GeV}^{-2}$ and is completely negligible at ordinary nuclear densities. However, at ultra-high density regimes these interactions become more important and, the neutral current interaction mediated by Z exchange acts coherently over the whole object, since only the weak charge matters and not the nature of the charge carriers [1]; so their effects must be analised.

The two-body potential arising from Z exchange in the static limit is obtained from the Fourier transform of the Z propagator and takes the Yukawa form [1] $V(r) = e^{-M_Z r}/4\pi r$, with an effective range of order $M_Z^{-1} \sim 10^{-2}$fm. At one loop, a long-range contribution from neutrino exchange modifies the potential [2, 3] with a correction that only dominates over the Yukawa potential only beyond a crossover distance $rM_z \simeq 15$. This long-range tail is cut off at distances of order $\sim 1\ \mu$m and, in the ultra-high density regime, the average distance between particles is many orders of magnitude smaller than this distance; therefore, the neutrino-mediated contribution is negligible, and the repulsive pressure analysed in this paper is provided due to the tree-level Z exchange [1].

For this weak interaction to act coherently, two conditions must be met simultaneously [1]: the weak charge density must trace the mass density, $\rho_W \propto \rho$, and the average separation between charge carriers, l, must satisfy $l \ll M_Z^{-1}$ to ensure that the charge distribution appears continuous at the scale of Z exchange. At the ultra-high density regime, these two conditions are satisfied, yielding a net repulsive interaction. It must be said that at such densities neutrons cannot exist, but since the net weak charge Q_W is conserved and the interaction does not depend on the nature of the carriers, the conclusion remains unchanged [1] as baryonic matter of u, d and s quarks is neutral and, from the point of view of its quantum numbers, behaves as a neutron star.

C. Equation of State at Ultra-High Densities

The initial analysis proposed by Zeldovich [6] states that the equation of state of the matter interacting by boson exchange can be derived from the energy per particle in the dense medium. For weak charge carriers of number density n, the pressure is,

$$p_W = \frac{G_F}{4\sqrt{2}} n^2 = D\ n^2, \tag{1}$$

where the number density of carriers is measured in fm^{-3} and the pressure in GeV^4 [1]. In this case, the energy density, equals $\varepsilon = \rho + p$, where ρ is the rest mass density. But in the scenario considered at ultra-high densities, $p \gg \rho$, therefore, the equation of state reduces to the stiff relation $\varepsilon \approx p$ which saturates the causality principle $dp/d\varepsilon = 1$ [6].

D. Degenerate Ultra-Relativistic Fermionic Gas

At extreme densities inside these compact stellar objects, the carriers of weak charge form a degenerate Fermi sea in which the Fermi momentum p_F exceeds the rest mass of the carriers. When we reach the ultra-relativistic limit, the rest mass contribution is negligible [11]. Therefore, the density of energy of the gas is obtained by integrating over all occupied momentum states up to the Fermi momentum p_F [11],$\varepsilon_{\rm deg} = \frac{g}{2\pi^2}\int_0^{p_F} p^3\,dp = \frac{g}{8\pi^2}p_F^4$, where g is the degeneracy factor accounting for spin and, in the case of quarks, colour degrees of freedom [9]. The number density of carriers is related to the Fermi momentum as [7], $p_F = \left(\frac{6\pi^2}{g}\right)^{1/3} n^{1/3}$, substituting this relation back to the density of energy [7, 11],

$$\varepsilon_{\rm deg} = \frac{g}{8\pi^2}\left(\frac{6\pi^2}{g}\right)^{4/3} n^{4/3} = \frac{3}{4}\left(\frac{6\pi^2}{g}\right)^{1/3} n^{4/3}. \tag{2}$$

If we consider the fermionic gas made of quarks, the degeneracy factor accounting for spin has to be computed as $g = g_c \times g_s$, where $g_c = 3$ accounts for the 3 colours that quarks have and $g_s = 2$ accounts for both spins. Therefore, $g = 6$, which are the indistinguishable degrees of freedom. We also have to consider that quarks possess 3 different flavours, which are distinguishable, so we will have a 3-factor multiplying the whole energy density. This brings to a final expression to,

$$\varepsilon_{deg} = \frac{9}{4}\pi^{2/3} \cdot n^{4/3} = K n^{4/3}, \tag{3}$$

with $K = \frac{9}{4}\pi^{2/3}$ The pressure of this gas follows the ultra-relativistic result $p = \varepsilon/3$ [7, 11]. In the expression 3 we must consider if QCD corrections are relevant for our problem. At extreme temperatures and densities characteristic of compact stellar objects, the strong coupling constant, α_s is sufficiently small due to asymptotic freedom, and considering that corrections to the degenerate fermionic gas are of order $\mathcal{O}(\alpha_s)$, we can therefore neglect these interaction terms [17].

E. The Tolman-Oppenheimer-Volkoff Equation

The structure of these relativistic compact objects follows the Tolman-Oppenheimer-Volkoff (TOV) equation [4, 5],

$$\frac{dp}{dr} = -\frac{G(\varepsilon + p)(M_\circ + 4\pi r^3 p)}{r(r - 2GM_\circ)}, \tag{4}$$

where the enclosed mass-energy in a radius r is defined, $M_\circ(r) = \int_0^r 4\pi r'^2 \varepsilon(r')\,dr'$. This equation is derived from Einstein's field equations,

$$R_{\mu\nu} = 8\pi G\left(T_{\mu\nu} - \frac{T}{2}g_{\mu\nu}\right), \tag{5}$$

applied to a static, spherically symmetric spacetime described by the Schwarzschild metric [12, 14], with a perfect fluid energy-momentum tensor [14], $T_{\mu\nu} = (\varepsilon + p)u_\mu u_\nu - p g_{\mu\nu}$, where ε is the total energy density, p the pressure, and u^μ the four-velocity of the fluid. It is important to note that the stellar objects predicted in [1] are extremely compact, with a Schwarzschild radius almost half of the physical radius. Therefore, to extend the analysis to the interior of the horizon, [1] employed Eddington-Finkelstein coordinates, in which the metric is non-singular at $r = r_s$, to show that the TOV equation is identical past the horizon, validating our use of the same equation for both sides of the horizon [1].

III. THEORETICAL FRAMEWORK AND SOLUTION METHOD

To analyse if the repulsive interactions arising from the Z-boson exchange are a relevant candidate for stabilizing a ultra-high density stellar object against gravitational collapse we have to consider a full equation of state that take into account the weak interaction pressure and the degeneracy pressure.

A. Equation of State

The total energy density of the system is the sum of the contributions due to the weak interactions and the degenerate Fermi sea. As we have seen, the equations of state regarding both effects are, respectively, $p_W \approx \varepsilon_W$ and $p_{deg} = \varepsilon_{deg}/3$, with the respective energy densities being (1), (3). Therefore, the total contribution to the density of energy amounts to the sum of the individual energy densities, while the contribution to the pressure takes into account the weight of the equation of state. In this case, our system is defined by the magnitudes:

$$\varepsilon = D\ n^2 + K\ n^{4/3}, \tag{6}$$

$$p = D\ n^2 + \frac{K}{3}\ n^{4/3}. \tag{7}$$

It is interesting to analyse the behaviour of the individual contributions to the total pressure according to the number density of carriers. To do so, we can consider the ratio p/ε in two different scenarios:

- When we are in the ultra-high density regime, computing $n \to \infty$ the ratio $p/\varepsilon \to 1$, therefore recovering the Zeldovich stiff condition [6].
- As we analyse the upper layers of the stellar object, density decreases rapidly, and by computing $n \to 0$ we obtain $p/\varepsilon \to 1/3$, recovering the pure ultra-relativistic gas result.

By considering these two extreme behaviours, we notice that the dominant mechanism that holds the object against gravitational collapse depends solely on the density of such an object; therefore, there must exist a crossover threshold where the dominant mechanism stops being the weak interaction repulsion and starts being the degenerate Fermi sea.

B. Crossover Density

The crossover density, n^*, marks the change in the dominant mechanism of structure in these stellar objects. Its relative position to the central density, n_c, and the surface density, n_R, informs us how much of the star is governed by each regime. If $n^* \ll n_R$ the entire star is in the weak interaction regime and the degeneracy correction is negligible, therefore recovering the results of [1]. If $n^* \gg n_c$ the entire star is degeneracy-dominated, and the weak interactions don't play any structural role. There is just one regime, whereas both mechanisms can play a non-negligible role when $n_r \lesssim n^* \lesssim n_c$, meaning that the core is weak-interaction-dominated and the outer layers are degeneracy-dominated.

To obtain a numerical value of this density, we must resolve when both contributions to the pressure are equal, $p_W = p_{deg}$. By solving this equation, the crossover density amounts to:

$$n^* = \left(\frac{3\sqrt{2}}{G_F}\right)^{3/2} \pi\ \text{GeV}^3 \approx 6.9 \cdot 10^8\ \text{GeV}^3 \tag{8}$$

Which roughly approximates to $n^* \approx 9 \cdot 10^{10}\ \text{fm}^{-3}$. With this result, by comparing it to the central and surface densities obtained through the TOV equation, we will be able to determine the relative importance of each structural mechanism.

C. TOV Equation with the Full Equation of State

As mentioned before, the structure of the star is determined through the TOV equation by imposing a certain equation of state. To characterise our system, we have to express the TOV equation (4) using the variables of the stellar object studied, which are the density of charge carriers, n, and the radius of the star, r. By imposing the equation of state defined in Section A, we can perform a change of variable in the relativistic equation of hydrostatic equilibrium that results in:

$$\left(2Dn + \frac{4K}{9}n^{1/3}\right)\frac{dn}{dr} = \frac{-G\left(2Dn^2 + \frac{4K}{3}n^{4/3}\right)}{r(r-2GM_o)}\left(M_o + 4\pi r^3\left(Dn^2 + \frac{K}{3}n^{4/3}\right)\right) \tag{9}$$

with $M_0 = \int_0^r 4\pi r^2 (Dn^2 + Kn^{4/3})dr$.

This integro-differential equation governs the physical characteristics of our system, and its initial conditions can be trivially determined by imposing a certain central density $n(0) = n_c$ and demanding that the enclosed mass at the centre is 0, $M_\circ(0) = 0$.

This equation can only be numerically integrated, as no analytical solution exist. To do such thing, we first have to notice the dimensional parameters that appear in our equation. The vast range of orders of magnitude that vary between these parameters would make it impossible to numerically integrate the equation without numerical instability or truncation errors; therefore, we must adimensionalise the TOV equation in order to integrate it [1].

D. Adimensionalization of the TOV Equation

To perform the adimensionalisation of the TOV equation, we must define adimensional variables from our initial ones, which are the number density and the radius. If we define :

$$y = \frac{K^{3/2}\sqrt{G}}{D}\ r \tag{10}$$

$$g(y) = \left(\frac{D}{K}\right)^{3/2} n(r(y)) \tag{11}$$

We can express the TOV equation through this adimensional variables which result in:

$$\begin{gathered}\left(2g + \frac{4}{9}g^{1/3}\right)\frac{dg}{dy} = \\ \frac{-\left(2g^2 + \frac{4}{3}g^{4/3}\right)}{y(y - 2\mu(y))}\left[\mu(y) + 4\pi y^3\left(g^2 + \frac{1}{3}g^{4/3}\right)\right]\end{gathered} \tag{12}$$

with the adimensional enclosed mass $\mu(y) = (K^{3/2}G^{3/2}/D)\ M_\circ(n(r(y)), r(y))$ being:

$$\mu(y) = \int_0^y 4\pi y^2 (g^2 + g^{4/3})\, dy. \tag{13}$$

Expressing Equation (13) as a differential equation,

$$\frac{d\mu(y)}{dy} = 4\pi y^2 (g^2 + g^{4/3}), \tag{14}$$

we notice that equations (12) and (14) form a system of coupled first-order Ordinary Differential Equations (ODE), therefore, we are facing an Initial Value Problem.

E. Numerical Integration

To integrate such a system of ODEs, we can use an adaptive Runge-Kutta integrator, for example, an 8th-order RK routine with the initial conditions $g(0) = g_c$ (adimensional central density) and $\mu(0) = 0$. By doing this, we first have to consider a few problems around the integration.

As we see in equation (12) the independent variable y presents a singularity at $y = 0$; therefore, we must start the integration at around $y = 10^{-6}$ to avoid such a thing. Also, in order to resolve the total enclosed mass, we must stop the integration at the surface of the stellar object, so we must define an event that marks such a boundary.

F. Boundary Condition

As stated before, to achieve the ultra-high density regimes necessary to even consider the weak interactions as a structural mechanism, we are analysing these stellar objects in the context of the primordial universe. For the purpose of this work, we shall assume that $T^* \sim 1$ GeV in order to analyse the stability of these objects, although further research should be done on the appropriate range of temperatures that ought to be used. In this case, the universe is composed of a thermal plasma of relativistic species that exert a certain radiation pressure [15] rendering invalid the classical definition of the surface of a stellar object following the TOV equation, $p(R) = 0$ [4, 5].It follows that the new boundary condition will arise from requiring hydrostatic equilibrium in the stellar boundary [7], $p_{interior}(n(r), T) = p_{exterior}(T)$. The pressure exerted by the plasma in the exterior of the star equals [15],

$$P_{exterior}(T) = \frac{\pi^2}{90} g_*(T) \cdot T^4 \tag{15}$$

with $g_*(T)$ being the relativistic degrees of freedom. Accounting for all relativistic species present at those temperatures the g-factor equals $g_* = 61.75$ [13], which turns Equation (23) in $p_{exterior}(T) = (247/360)\pi^2\ T^4$. At the temperature considered, $T^* \sim 1$ GeV, the pressure in the exterior becomes $p_{exterior}(T^*) = (247/360)\pi^2 \mathrm{GeV}^4$. At surface level, the only relevant contribution to the pressure is going to be due to the degenerate Fermi sea, $p_{interior} = p_{deg}$, as densities reduce and weak interactions become negligible. In future work, it must be analysed if thermal corrections must be added to expression 2, but considering a simple model without those, our boundary condition equals to:

$$n(R) = \left(\frac{247}{270}\right)^{3/4} \pi \approx 382.5\ \mathrm{fm}^{-3}, \tag{16}$$

which is more than 2000 times more than the nuclear saturation density. This density in adimensional units equals $g_{surface} \approx 8.7 \cdot 10^{-9}$; therefore, we must stop the numerical integration when the value of the adimensional density reaches the value of $g_{surface}$.

IV. RESULTS

A. Density and Mass Profiles

By performing the numerical integration starting at $y = 10^{-6}$ with the boundary condition set to $g_{surface}$, we obtain the following density profiles for different initial values of the central density g_c:

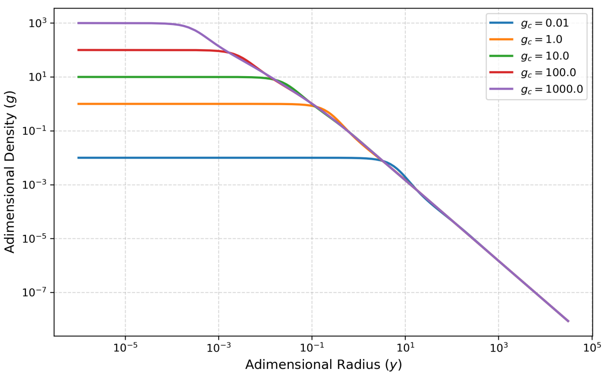

Figure 1: Adimensional density profiles in logarithmic scale for different central density values

These density profiles show a normal behaviour for a stellar object of these characteristics. In the centre, density appears constant, while it decreases as the radius augments. The interesting behaviour that Figure 1 presents is the sudden convergence of all the profiles to the same point, which means that independently of their central density, all the stellar objects meet the boundary condition at around the same radius, $y_{max} \approx 30870$, which in SI units equals $R \approx 15$ m. This universal behaviour is a consequence of the polytropic degeneracy pressure $p \propto n^{4/3}$ dominating the envelope of the star at low densities, which acts as an attractor solution in the TOV equation independent of the initial conditions [7, 11].

To analyse the stability of these solutions, we have to make use of the turning point theorem of stellar stability [7, 16] where it is stated that a sequence of stable configurations becomes unstable at the maximum of the curve $M_\circ(n_c)$ beyond which no stable equilibrium exists [7]. To do so, we can compute the enclosed mass inside the radius y_{max} as a function of the central density of our stellar object, g_c:

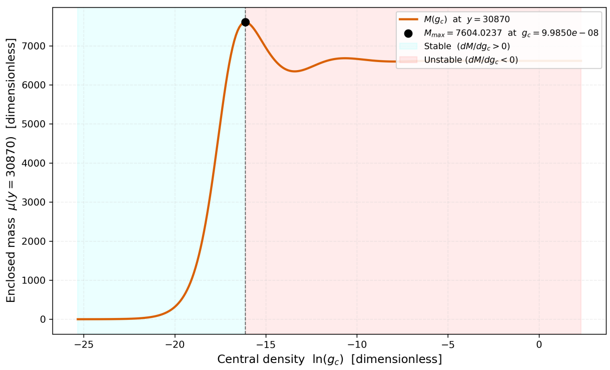

Figure 2: Dimensionless enclosed mass at $y_{max} = 30870$ as a function of the central density g_c

This maximum is found at $g_c^{max} = 9.98 \cdot 10^{-8}$ with a height of $\mu_{max}(y = 30870) \approx 7604$, which, if we dimensionalise using the inverted relation in Equation (11), amounts to approximately $n_c^{max} \approx 4387.7$ fm^{-3} and $M_\circ^{max} \approx 2.4 \cdot 10^{-3}\ M_\odot$ (in solar masses).

This result becomes important when we compare it to the crossover density that we have previously defined. With $n^* \sim 10^{10}$ fm^{-3} and $n_c^{max} \sim 10^3$ fm^{-3} we clearly see that $n^* \gg n_c^{max}$, therefore implying that all stable configurations of our stellar object are entirely degeneracy-dominated, which leads us to the conclusion that weak interactions are not a relevant structural mechanism in our system.

V. CONCLUSIONS

- Considering both the weak interaction and degenerate Fermi sea contributions to the total internal energy defines a crossover density, $n^* \approx 9 \cdot 10^{10}$ fm^{-3}, where, by comparing it to the density profile of the stellar object, the relative importance of each contribution can be assessed.
- By defining a boundary condition such as $p_{interior} = p_{exterior}$ a universal radius of $R \approx 15$ m is found independently of the initial central density of the object.
- By analysing the stability of the solutions obtained, we have determined that the central densities required for weak interactions to be important structural mechanisms cannot be reached while in the stable branch, therefore, rendering such interactions irrelevant for this problem.
- The evolution of these compact objects can be a further branch of research, as, after the universe cools down, such objects would intuitively expand without growing in mass, leading to the formation of ultra-light neutron stars, whose creation mechanism are still unknown.

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Interaccions febles com a mecanisme de suport contra el col · lapse gravitatòri en estrelles bariòniques primordials.

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Resum: En aquest treball s'ha investigat si les interaccions febles mediades per l'intercanvi de bosons Z poden aportar una pressió estabilitzadora conra el col · lapse gravitatòria en estrelles de bariònics primordials. Desenvolupant el treball d'Espriu , s'ha completat l'equació d'estat incorporant-hi la pressió de degeneració d'un gas fermiònic ultrarelativista. S'ha definit una densitat de traspàs , $n^* \approx 1,9 \cdot 10^9$ fm^{-3}, que marca la transició entre el règim dominat per és interaccions febles i per la degeneració. S'ha resolt l'Equació de Tolman-Oppenheimer-Volkoff numèricament amb una condició de frontera derivada de l'equilibri hidroestàtic amb la pressió de radiació tèrmica del plasma primordial, aconseguint la deducció d'un radi universal d'aproximadament 9 metres. A través d'una anàlisi d'estabilitat usant el "Turning Point theorem" s'ha establert que totes les configuracions estables posseeixen una densitat central menor a la densitat de traspàs, demostrant que les interaccions febles no juguen un paper rellevant en l'estructura d'aquests objectes compactes.
Paraules clau: Estrella Bariònica Primordial, Interaccions Febles, Gas Degenerat de Fermi, Equació de Tolman-Oppenheimer-Volkoff.
ODSs: Aquest TFG està relacionat amb els Objectius de Desenvolupament Sostenible 4 i 17.

Objectius de Desenvolupament Sostenible (ODSs o SDGs)

1. Fi de la es desigualtats		10. Reducció de les desigualtats	
2. Fam zero		11. Ciutats i comunitats sostenibles	
3. Salut i benestar		12. Consum i producció responsables	
4. Educació de qualitat	x	13. Acció climàtica	
5. Igualtat de gènere		14. Vida submarina	
6. Aigua neta i sanejament		15. Vida terrestre	
7. Energia neta i sostenible		16. Pau, justícia i institucions sòlides	
8. Treball digne i creixement econòmic		17. Aliança pels objectius	X
9. Indústria, innovació, infraestructures			