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It is shown that for all stars for which the radiation-pressure is greater than a tenth of the total pressure, an appeal to the Fermi-Dirac statistics to avoid the central singularity which arises in the discussions of the centrally condensed and the collapsed stars cannot be made. The bearing of this result on the possible state of matter in the interior of stars is indicated.

Since the publication of Milne’s memoir on the “Analysis of Stellar Structure” in the Monthly Notices for November 19301) a great deal of work has been done to consider “composite” stellar models. But the following simple considerations seem to have escaped notice and it seems worth while to state them explicitly.

§ 1. The Surfaces of Demarcation. As we approach the centre of a centrally-condensed or a collapsed star, we change over to the equation of state \( p = K_1 \varepsilon^{1/4} \) if the perfect gas law breaks down. If the perfect gas-law breaks down at all, the actual transition from the perfect-gas envelope to the degenerate core must occupy a certain zone, but we could for the sake of convenience consider a definite surface of demarcation defined as the surface at which the two equations of state give the same gas-pressure.

Now in the perfect gas envelope the total pressure is given by

\[
P = \left[ \left( \frac{k}{\mu} \right)^{1/4} \frac{\beta}{a} \left( \frac{1-\beta}{\beta} \right)^{1/8} \right] \varepsilon^{4/3},
\]

where

\[
\beta = 1 - \frac{\kappa L}{4\pi c G M},
\]

\(\kappa\) = the opacity coefficient, \(L\) = luminosity in ergs cm\(^{-2}\) s\(^{-1}\), \(M\) = mass in grams, \(k\) = Boltzmann’s Constant, \(\mu\) = molecular weight = \(\alpha m_H\) (say), \(m_H\) = mass of the hydrogen atom. Since for the standard model the gas pressure \(p\) is given by

\[
p = \beta P
\]

we have

\[
p = C \varepsilon^{4/3},
\]

where

\[
C = \left[ \left( \frac{k}{\mu} \right)^{1/4} \frac{1-\beta}{\beta} \right]^{1/8} \frac{2.632 \cdot 10^{15}}{\alpha^{4/3}} \left( \frac{1-\beta}{\beta} \right)^{1/8}.
\]

1) Refered to as l.c.
The equation of state in the degenerate zone is
\[ p = K_1 \varepsilon_{1a}^{5/3} , \] (5)
where
\[ K_1 = \frac{1}{20} \left( \frac{8}{\pi} \right)^{2/3} \frac{\hbar^3}{m^2 \mu_{1/3}^{1/3}} = \frac{9.890 \cdot 10^{12}}{\alpha^{5/3}} . \] (6)

At the first surface of demarcation which we will call \( S_1 \), the density \( \varepsilon_1 \) is given by
\[ C \varepsilon_1^{1/3} = K_1 \varepsilon_{1a}^{5/3} , \]
or
\[ \varepsilon_1 = \left( \frac{C}{K_1} \right)^{3} . \] (7)

Now, it is well known that the equation of state (5) changes over again into the relativistic-degenerate-equation of state
\[ p = K_2 \varepsilon_{1a}^{4/3} , \] (8)
where
\[ K_2 = \frac{\hbar c}{8 \mu_{1/3}^{1/3}} \left( \frac{3}{\pi} \right)^{1/3} = \frac{1.228 \cdot 10^{15}}{\alpha^{4/3}} . \] (8')

Hence if circumstances permit we have to consider a second surface of demarcation, \( S_2 \), where the density \( \varepsilon_2 \) is given by
\[ \varepsilon_2 = \left( \frac{K_2}{K_1} \right)^{3} . \] (9)

Hence we have two surfaces of demarcation if and only if
\[ \varepsilon_2 > \varepsilon_1 , \]
or
\[ \left( \frac{K_2}{K_1} \right)^{3} > \left( \frac{C}{K_1} \right)^{3} . \] (10)

i.e. only when [cf. equations (4'), (6), (8')] \[ \frac{1 - \beta}{\beta} < \frac{\hbar^3 c^5 a}{512 \pi l^4} = 0.1015 , \]
or
\[ \beta > 0.9079 . \] (11)

It may be remarked in passing that the above value for \( \beta \) is independent of the assumed molecular weight. It depends only on the mass, luminosity and opacity in the gaseous envelope. It is also independent of whether we consider the same opacity for the degenerate zone and the gaseous envelope, or different opacities in the two regions.
§ 2. The meaning of the fundamental inequality (11) is made clear by the following.

In the following graph I plot \( \log p \) against \( \log \varrho \).

For numerical calculations I use \( \alpha = 2 \). The straight line \( ABK \) represents the equation of state \( p = K_1 \varrho^5 \) and \( BC \) the equation of state \( p = K_2 \varrho^4 \). These two intersect at \( B \) where the density is that which corresponds to the second surface of demarcation, namely \( \varrho_2 \). \( ABC \) gives roughly the equation of state of a degenerate gas.

Let us consider a star for which \( \beta = 0.98 \). By (4) we get

\[
\log p = 14.455 + \frac{4}{3} \log \varrho.
\]  

(12)

\( DE \) represents this equation. It intersects the degenerate equation of state \( A, B, C \) at \( E \). The point \( E \) corresponds to the first surface of demarcation \( S_1 \). Hence for all stars for which \( \beta = 0.98 \), we first traverse a perfect gas envelope with an equation of state represented by \( DE \). Then we traverse a degenerate zone corresponding to \( EB \) and finally (if we have not yet reached the centre) a relativistically degenerate zone.

Now, if \( \beta = 0.9079 \), then \( GB \) represents the perfect gas equation of state and the degenerate zone reduces to a single layer, and the relativistically degenerate zone is described equally well by the perfect gas equation.

Now if \( \beta < 0.9079 \) the perfect gas equation of state has no intersections with \( ABC \) and this means that however high the density may become the temperature rises sufficiently rapidly to prevent the matter from becoming degenerate.

In this connection it will have to be remembered that considerations of relativity do not affect the equation of state of a perfect gas. \( p = N k T \), is true independent of relativity.

§ 3. Centrally-Condensed Stars. Now, for each mass \( M \) there is a unique luminosity \( L_0 \) — the "Eddington luminosity" which makes the star

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1) The radiation pressure is greater than a tenth of the total pressure if \( \beta < 0.9079 \).
a perfect gas sphere, with a polytropic index 3. This \( L_0 \) characterizes a unique \( \beta_0 \) which is in fact related to \( M \) by means of Eddington's quartic equation:

\[
1 - \beta = 0.00809 \left( \frac{M}{\odot} \right)^{\alpha} \beta^4.
\]  

(13)

Now from the definition of a centrally-condensed and a collapsed star, it is clear that

\[
\beta_\text{c-c} < \beta_0 \quad | \quad \beta_\text{col.} > \beta_0.
\]  

(14)

Consider first the mass \( M \) for which \( \beta_0 = 0.9079 \). By (13) we have

\[
\frac{M}{\odot} = 6,623 \alpha^{-2}.
\]  

(15)

If we assume \( \alpha = 2 \),

\[
\frac{M}{\odot} = 1.556.
\]  

(15')

Now consider a centrally-condensed star of mass \( M \) greater than (or equal to) \( M \). Then we obviously have

\[
M \beta_0 < \frac{\odot}{M} \beta_0 = 0.908.
\]

\[
M \beta_\text{c-c} < \frac{\odot}{M} \beta_0 < 0.908.
\]  

(16)

Hence, we have the result that for all centrally condensed stars of mass greater than \( M \), the perfect gas equation of state does not break down, however high the density may become, and the matter does not become degenerate. An appeal to the Fermi-Dirac statistics to avoid the central singularity cannot be made.

Since however we cannot allow the infinite density which the centrally condensed solution of Emden's differential equation — index 3 — allows at the centre and in the absence of our knowledge of any equation of state governing the perfect gas other than that of degenerate matter, our only way out of the singularity is to assume that there exists a maximum density \( \rho_\text{max} \) which matter is capable of. We have therefore to consider the "fit" of a gaseous envelope of the centrally condensed type on to a homogeneous core at the maximum density of matter. If we insist on the density to be continuous at the interface the equation of "fit" is found to be

\[
\frac{1}{3} \xi' \Theta'^3 = -\left( \frac{d \Theta}{d \xi} \right)_{\xi = \xi'}.
\]  

(17)

where the polytropic equation describing the gaseous part of the star is

$$\frac{1}{\xi^2} \frac{d}{d \xi} \left( \xi^2 \frac{d \Theta}{d \xi} \right) = -\Theta^3,$$  

(17')

where $\xi'$ is the value of $\xi$ at which $\Theta_{\text{max}}$ begins. In (17') the meaning of $\Theta$ and $\xi$ are the following:

$$\Theta = \lambda_3 \xi^3, \quad \xi = \frac{C}{\pi G \beta} \lambda_3^{-1/3}. \tag{17''}$$

($\lambda_3$ is a homology constant). But (17) has no solutions if $\Theta$ is of the Emden's or of the centrally-condensed type. Hence the acceptance of a $\Theta_{\text{max}}$ does not help us out of the difficulty if we insist on the density to be continuous at the interface. The procedure then to construct an equilibrium configuration would be to proceed along the centrally condensed solution until the mean density $\bar{\rho}(r)$ of the surviving mass $M(r)$ equals $\Theta_{\text{max}}$ which will occur at a determinate $r = r''$ (say) where

$$M(r'') = \frac{4}{3} \pi r''^3 \Theta_{\text{max}}; \tag{18}$$

we then replace the material inside $r = r''$ by a sphere of incompressible matter at the density $\Theta_{\text{max}}$. At $r''$ there will be a discontinuity of density (see Fig. 2).

Now the form of $\Theta$ as $\xi \to 0$ for a centrally condensed solution is (Milne, l. c.):

$$\Theta \sim \frac{1/2}{\xi \left[ \log \left( D/\xi \right) \right]^{1/2}}, \tag{19}$$

where $D$ is a constant. $D$ is fixed by the condition that the analytic continuation of (19) passes through $\xi = 1$ and $\Theta = 0$ and satisfies here the requisite boundary condition, namely

$$M = -\frac{4}{\pi^{1/2} \left( C \beta \right)^{3/2}} \left( \frac{\xi^2 d \Theta}{d \xi} \right). \tag{20'}$$

1) $C$ is given by equation (14').
Hence we get the result that $D$ is a function of $L$, $M$ and $x$ only and hence fixed. Since $D$ is fixed by the boundary condition, it follows that the value of $\xi''$ at which $\Theta'(\xi'')$ becomes equal to $\Theta_{\text{max}}$ (where cf. equation (17''))

$$\Theta_{\text{max}} = \frac{1}{4} \frac{1}{3} \frac{1}{3},$$

(21)

is fixed as a function of $\lambda_g$. In other words the discontinuity in $\Theta$, $\Delta \Theta''$ at the interface $\xi''$ is a single-valued function of $L$, $M$, $x$ and $\lambda$ or

$$\Delta \Theta'' = F(L, M, x; \lambda_g)$$

(21')
or by (21)

$$\Delta \varrho'' = f(L, M, x; \lambda_g)$$

(22)

where $\Delta \varrho''$ is the discontinuity of density at the interface.

But it has been suggested by Landau (1) (among others) that the maximum density of matter will arise after some kind of overcompressibility, the incompressibility setting in later (see Fig. 3).

Further it has been suggested that 1) the pressure at which the overcompressibility sets in must be a physical property of the atomic nuclei and the electrons in the enclosure, and 2) the form of the curve $ABC$ is again an intrinsic physical property of matter. If we idealise the situation of Fig. 8, we see that $\Delta \varrho$ ought to be a physical property of matter. Let this $\Delta \varrho$ be $\omega$. Then by (22) we have to so choose the homology constant $\lambda_g$, that $\Delta \varrho''$ equals $\omega$:

$$f(L, M, x; \lambda_g) = \omega.$$ 

(23)

This fixes $\lambda_g$ and hence by (17'') fixes $r_0$ — the radius of the configuration. Hence we are able to obtain equilibrium configurations for arbitrary mass, and arbitrary luminosity, the radius however being determinate in each case.

§ 4. In the above section we have tried to construct the equilibrium configurations for all centrally-condensed stars of mass greater than $M^2$, and found that the introduction of a homogeneous core at the maximum density of matter ($\varrho_{\text{max}}$) with a discontinuity of density at the interface was necessary. We may now ask about the equilibrium configurations for centrally condensed stars with $\beta > 0.908$. Now the star has clearly a degenerate zone (see Fig. 1). A little consideration shows that if we come along a centrally-condensed solution in the perfect gas part of the star then at the interface $S_1$ (cf. § 1) we are compelled to choose a centrally-condensed solution for the polytropic equation of index $"{3/2}"$

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1) I am indebted to Dr. Strömgren for advice on these matters.

2) Or more generally, centrally-condensed stars with $\beta < 0.908$. 

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to describe the non-relativistic degenerate part of the star\(^1\)); also at the
second surface demarcation \(S_2\) we are again forced to choose a centrally-
condensed solution for the polytropic equation of index "3". Hence in
this case also we are unable to avoid the central singularity by appeal-
ing to the Fermi-Dirac statistics alone. The star must have a homo-
genous core with a discontinuity of density at the interface. The con-
siderations of the previous section apply and we see that the centrally-
condensed stars \(\beta > 0.908\) differ from the centrally-condensed stars with
\(\beta < 0.908\) only in this, that while in the former type of stars we have
to traverse a degenerate zone before reaching the homogeneous core, in
the latter type, the stellar material continues to be a perfect gas till we
reach the homogeneous core. Thus we find that all centrally-condensed
stars (on the standard model) must have a homogeneous core at the centre
with a discontinuity of density at the interface.

§ 5. Collapsed-Stars. Just a few remarks about collapsed stars may
be permitted. A detailed analysis of highly-collapsed stars has been given
elsewhere (Chandrasekhar, l.c.).

Consider a collapsed star of mass greater than \(\mathcal{M}\) and let further
\(\beta_0 < \beta_{\text{col}} < 0.9078\). In other words the "collapse" has not proceeded
sufficiently far to increase \(\beta\) beyond 0.9078. In such a case the collapse
can occur only on a homogeneous core. But if the collapse proceeds
sufficiently far, such that \(\beta_{\text{col}} > 0.9078\) in spite of \(\beta_0\) being less than
0,9078, the star will then possess a degenerate zone as well.

Conclusion: We may conclude that great progress in the analysis
of stellar structure is not possible before we can answer the following
fundamental question:

Given an enclosure containing electrons and atomic nuclei, (total charge
zero) what happens if we go on compressing the material indefinitely?

\(^1\) This is also true if we ascribe different opacities to the gaseous
and the degenerate part of the star.