Effects of Strange Particles on Neutron-Star Cores

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We show that uncertainties in the strength of interactions of hyperons among themselves and with nucleons lead to an uncertainty in the maximum allowed neutron-star mass of nearly a factor of 2, even if the properties of nuclear and neutron matter were known with infinite precision around normal nuclear density and below. The possibility of a transition to quark matter places some constraint on the strength of the hyperonic interactions.

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There has been tremendous interest and activity for nearly three decades in the relation between the nuclear-matter equation of state and the properties of neutron stars. This interest has intensified in the last decade with the theoretical activity in the simulation of the supernova event itself,[1] and with the activity surrounding heavy-ion collisions and efforts to extract nuclear-matter properties from them.[2] Much recent emphasis has focused on the stiffness of nuclear matter, particularly in the density range $\frac{1}{2} \lesssim n/n_0 \lesssim 2$ ($n_0 = 0.153$ fm^{-3}), although certainly the compressibility K is a density-dependent quantity.[3,4] At higher densities (usually $n/n_0 > 2$ but this is model dependent), hyperons appear due to strangeness-changing weak interactions. Although hyperons have been included in many (but not all) of the commonly used equations of state[5] a systematic study is still lacking, as emphasized recently.[4] Here we report on the somewhat surprising sensitivity of neutron-star masses to the uncertainties in the strength of hyperon interactions among themselves and with nucleons. It is a real challenge to nuclear physics to determine these interactions with more precision than is presently available.

We use the theoretical framework of relativistic nuclear mean-field theory.[6] Other theoretical frameworks are possible: The advantage of this one is that the resulting equation of state is automatically consistent with special relativity and it readily incorporates hyperons and their interactions. In our opinion it is best to view relativistic nuclear mean-field theory as a convenient way of parametrizing the equation of state. Therefore the parameters involved may differ somewhat from their values determined on some other basis, for example, phase-shift analyses of nucleon-nucleon scattering.

In our calculations we allow for the presence of n, p, Λ, Σ^-, e^-, and μ^- in the star. Strong interactions are mediated by the mean fields of the mesons ω, ρ^-, and a scalar σ which are generated by the presence of the baryons. The pressure is

$$P = \tfrac{1}{2} m_\omega^2 \bar{\omega}^2 + \tfrac{1}{2} m_\rho^2 \bar{\rho}^2 - \tfrac{1}{2} m_\sigma^2 \bar{\sigma}^2 + \bar{U}(\sigma) + \sum_i P_{\rm FG}(\bar{\mu}_i, \bar{m}_i) , \qquad (1)$$

where

$$\begin{aligned} g_\omega \bar{\omega} &= (g_\omega/m_\omega)^2 [n_p + n_n + R_\omega (n_\Sigma + n_\Lambda)] , \\ g_\rho \bar{\rho} &= (g_\rho/m_\rho)^2 [\tfrac{1}{2}(n_p - n_n) - R_\rho n_\Sigma] , \\ g_\sigma \bar{\sigma} &= (g_\sigma/m_\sigma)^2 [n_p^s + n_n^s + R_\sigma (n_\Sigma^s + n_\Lambda^s) \\ &\qquad - b m_N (g_\sigma \bar{\sigma})^2 - c (g_\sigma \bar{\sigma})^3] \end{aligned} \qquad (2)$$

are the mean fields. The coupling of ω to nucleons is denoted by g_ω, its coupling to hyperons relative to nucleons by $R_\omega = g_{\omega\Lambda}/g_\omega = g_{\Sigma\omega}/g_\omega$, and similarly for the other meson couplings. The particle densities are related to the Fermi momenta by $n_i = k_i^3/3\pi^2$ and the n_i^s are scalar ($\langle \bar{\psi}\psi \rangle$) densities.[6] The effective chemical potentials $\bar{\mu}_i = (\bar{m}_i^2 + k_i^2)^{1/2}$ are related to the true chemical potentials by

$$\begin{aligned} &\mu_n = \bar{\mu}_n + g_\omega \bar{\omega} - \tfrac{1}{2} g_\rho \bar{\rho} , \quad \mu_p = \bar{\mu}_p + g_\omega \bar{\omega} + \tfrac{1}{2} g_\rho \bar{\rho} , \\ &\mu_\Lambda = \bar{\mu}_\Lambda + R_\omega g_\omega \bar{\omega} , \quad \mu_\Sigma = \bar{\mu}_\Sigma + R_\omega g_\omega \bar{\omega} - R_\rho g_\rho \bar{\rho} \\ &\mu_e = \bar{\mu}_e , \quad \mu_\mu = \bar{\mu}_\mu . \end{aligned} \qquad (3)$$

The baryons have effective masses

$$\begin{aligned} &\bar{m}_n = m_n - g_\sigma \bar{\sigma} , \quad \bar{m}_p = m_p - g_\sigma \bar{\sigma} , \\ &\bar{m}_\Lambda = m_\Lambda - R_\sigma g_\sigma \bar{\sigma} , \quad \bar{m}_\Sigma = m_\Sigma - R_\sigma g_\sigma \bar{\sigma} . \end{aligned} \qquad (4)$$

The chemical potentials must obey the equations of chemical equilibrium and electrical neutrality. There is only one independent chemical potential, which is the baryon chemical potential μ_n. The nonlinear potential is $U(\bar{\sigma}) = -\frac{1}{3} b m_N (g_\sigma \bar{\sigma})^3 - \frac{1}{4} c (g_\sigma \bar{\sigma})^4$. Finally $P_{\rm FG}$ is the relativistic Fermi-gas pressure.

Apart from the hyperon couplings this equation of state has five parameters: g_σ/m_σ, g_ω/m_ω, g_ρ/m_ρ, b, and c. These are determined by four macroscopic properties of isospin-symmetric nuclear matter at saturation: density[7] $n_0 = 0.153$ fm^{-3}, binding energy[7] 16.3 MeV, symmetry-energy coefficient[7] $a_{\rm sym} = 32.5$ MeV, and compressibility[7,8] $K = 300$ MeV. The fifth parameter is determined by a microscopic property, the Landau mass[9] $m_L = (m_N^{-2} + k_F^2)^{1/2} = 0.83 m_N$. Since uncertainties in these numbers have been considered elsewhere[3,4] we simply fix and focus on the hyperons.

The strengths of the couplings of hyperons relative to nucleons, R_σ, R_ω, R_ρ, are of immediate concern. First consider the complete omission of hyperons. It is known[10] that relativistic nuclear mean-field theory yields an equation of state which is nearly identical to the highly regarded one of Friedman and Pandharipande,[11] which is based on two-body free-space potentials and phenomenological three-body forces, when K and m_L are adjusted to match. Solving the Oppenheimer-Volkoff-Tolman equation yields the neutron-star mass as a function of central density as labeled by N in Fig. 1. The maximum mass is $2.1M_\odot$. At the other extreme we could allow the presence of hyperons but decouple them from the mean meson fields, $R_\sigma = R_\omega = R_\rho = 0$. The resulting mass curve is labeled 0 in Fig. 1. The maximum mass is reduced to about $1.15M_\odot$. In going from N to 0, nucleons are converted to hyperons with increasing density. This reduces the energy stored in the repulsive vector fields and softens the equation of state. Neither of these extremes seems likely.

In his studies Glendenning[5] assumed universality among the coupling constants, $R_\sigma = R_\omega = R_\rho = 1$. This is not unreasonable in the absence of any information. But Boguta and Bohrmann[12] have studied the energy levels of Λ hypernuclei within the framework of relativistic mean-field theory. Using the values $R_\sigma = R_\omega = \frac{1}{3}$ (equality of R_σ and R_ω is suggested by the quark model) they obtained a good representation of the available data. When more data became available, Rufa *et al.*[13] made a reanalysis. With the constraint $R_\sigma = R_\omega$ they obtained 0.21 ± 0.02. When that constraint was relaxed it was found that $\frac{1}{2}(R_\sigma + R_\omega) = 0.52 \pm 0.48$, $R_\omega - R_\sigma = 0.08 \pm 0.13$, a tremendous uncertainty (the $\chi^2/N_{\rm DF}$ are 4.4 and 3.6, respectively). In the absence of further information we assume that $R_\rho = R_\omega$, and that Σ's behave identically to Λ's apart from their vacuum masses. The resulting mass curves are plotted in Fig. 1 (in one case, labeled 0.48, we have taken $R_\omega = 0.56$ and $R_\sigma = 0.48$). There is a large uncertainty induced in the maximum mass by uncertainties in the strengths of the hyperon couplings.

The large spread in the values of the maximum mass seen in Fig. 1 can be compared to known masses of neutron stars. A good example is PSR 1913+16, whose mass has been determined to be[14] $M = (1.442 \pm 0.003)M_\odot$. Our equation of state is only compatible with this observation for $R_\sigma \gtrsim 0.4$.

More uncertainties may be caused by effects not considered here. The Ξ^- will appear at the higher densities. This would further soften the equation of state. Since it has strangeness -2 its couplings are even more remotely related to those of the nucleons. In addition, a ϕ-meson condensate may be generated by the presence of hyperons. Being a vector meson it would stiffen the equation of state. However, in the naive quark model, its presence would not affect the curves labeled N and 0 but would only shift the interior curves. These topics will be taken up in a more detailed paper.

Is there a phase transition to quark matter before gravitational instability sets in? We describe the quark phase in terms of the leptons, massless u and d quarks, a massive s quark ($m_s = 180$ MeV), one-gluon-exchange interactions among the quarks,[15] and a bag constant

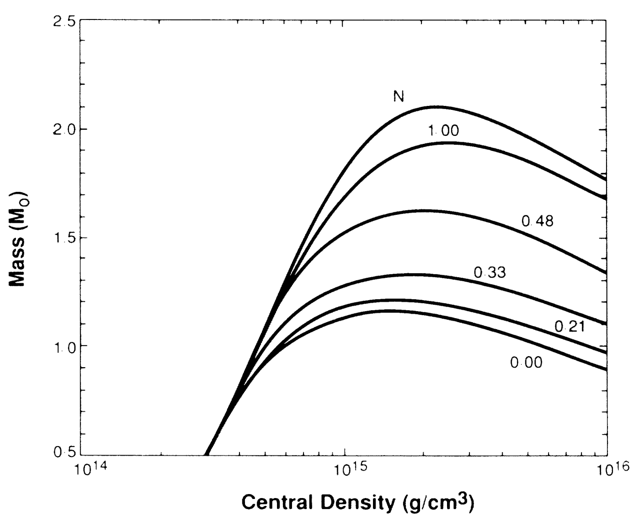

FIG. 1. The masses of neutron stars as functions of the central density. The curves display the sensitivity to hyperon interactions. The curve labeled N does not include hyperons. The other curves are labeled by the value of $R_\sigma = R_\omega$ (with the exception of the curve labeled $R_\sigma = 0.48$ in which case $R_\omega = 0.56$) used in the equation of state. The curve $R_\sigma = 0$ represents a case where hyperons are present as noninteracting states.

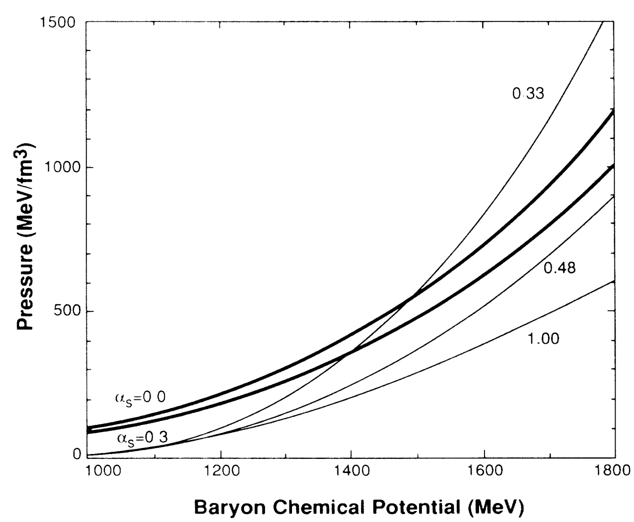

FIG. 2. The equation of state for nuclear matter (thin lines labeled by $R_\sigma = 0.33$, 0.48, and 1.00) and for quark matter (thick lines labeled by $\alpha_s = 0.0$ and 0.3). Not all choices of R_σ and α_s are thermodynamically consistent with a phase transition to quark matter.

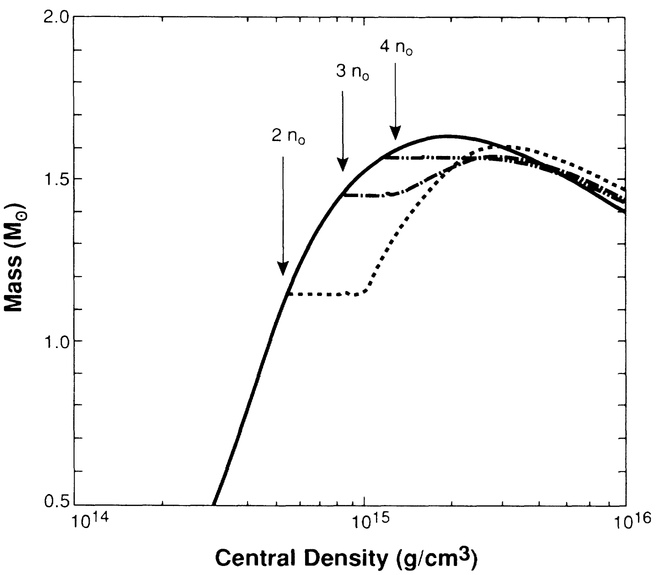

FIG. 3. The masses of neutron stars assuming a phase transition to quark matter at $n_B = 2n_0$, $3n_0$, and $4n_0$ as labeled. The nuclear equation of state assumes $R_\sigma = 0.48$. The horizontal portion of the broken curves represents the density discontinuity associated with the phase transition.

$B > 0$ to simulate confinement. The kinetic and interaction pressure (neglecting the additive bag term $-B$) for several values of the strong coupling α_s is plotted against the baryon chemical potential in Fig. 2, as are several of the nuclear pressures. Reasonable values of B lie in the range from 50 to 450 MeV/fm^3. It is clear from this figure that the R's and α_s cannot be chosen arbitrarily, but their allowed values must be correlated. For example, $R = \frac{1}{3}$ would imply that even at very high density nuclear matter would be thermodynamically favored. When $R_\sigma = 0.48$ the nuclear curve would be inconsistent with $\alpha_s = 0.3$ because, with suitably chosen but not atypical values of B, quark matter would be stable at both low and high densities whereas nuclear matter would be stable at intermediate densities. However, $\alpha_s = 0.2$ poses no such inconsistencies. Since the underlying dynamics is QCD it should be no surprise that α_s and the R's must be correlated, although in practice the relation cannot be presently determined.

The mass curves of Fig. 3 show that a quark core is possible if $R_\sigma = 0.48$, $\alpha_s = 0.2$, and if the transition out of the nuclear phase begins at $2n_0$ or $3n_0$. If it begins at $4n_0$ gravitational instability sets in and no quark core is possible. Further parameter choices will be considered in an upcoming paper.

We can estimate the maximum Keplerian frequency of rotation from the formula

$$\Omega_K = (7.2\times 10^3\ \mathrm{s}^{-1})(M_s/M_\odot)^{1/2}[R_s/(10\ \mathrm{km})]^{-3/2}, \quad (5)$$

where M_s and R_s are the maximum mass and corresponding radius for the spherical nonrotating star.[16] Fixing $R_\sigma = 0.48$, we find that in the absence of a transition to quark matter $\Omega_K = 7400\ \mathrm{s}^{-1}$, whereas with a transition at $n = (2\text{–}3)n_0$, $\Omega_K = 9700\text{–}9100\ \mathrm{s}^{-1}$. In these cases we cannot obtain a frequency as high as that reported[17] for the remnant of SN 1987A, $\Omega = 12\,370\ \mathrm{s}^{-1}$.

Our main conclusion is that even if we had perfect knowledge of the nuclear equation of state up to about $2n_0$ there is still a large uncertainty in the maximum neutron-star mass induced by uncertainties in the strengths of the hyperon interactions. It is a very real challenge to nuclear physics to determine these strengths in laboratory experiments. Our other conclusion is that consistency between the nuclear- and quark-matter equations of state must be maintained. There are probably commonly used nuclear equations of state which are incompatible with a transition to quark matter. Both of these issues should be investigated in theoretical frameworks other than relativistic nuclear mean-field theory.

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[1] *Theory of Supernovae*, edited by G. E. Brown [Phys. Rep. **163**, 1 (1988)].

[2] J. Hüfner, Phys. Rep. **125**, 129 (1985); L. P. Csernai and J. I. Kapusta, *ibid.* **131**, 223 (1986); R. Stock, *ibid.* **135**, 259 (1986); H. Stöcker and W. Greiner, *ibid.* **137**, 277 (1986); R. B. Clare and D. Strottman, *ibid.* **141**, 177 (1986).

[3] N. K. Glendenning, Phys. Rev. Lett. **57**, 1120 (1986).

[4] M. Prakash, T. L. Ainsworth, and J. M. Lattimer, Phys. Rev. Lett. **61**, 2518 (1988).

[5] V. A. Ambartsumyan and G. S. Saakyan, Astron. Zh. **37**, 193 (1960) [Sov. Astron. **4**, 187 (1960)]; V. R. Pandharipande, Nucl. Phys. **A178**, 123 (1971); H. A. Bethe and M. Johnson, *ibid.*, **A230**, 1 (1974); N. K. Glendenning, Z. Phys. A **327**, 295 (1987).

[6] B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. **16**, 1 (1986).

[7] P. Möller, W. D. Myers, W. J. Swiatecki, and J. Treiner, At. Data. Nucl. Data Tables **39**, 225 (1988).

[8] M. M. Sharma *et al.*, Phys. Rev. C **38**, 2562 (1988).

[9] C. H. Johnson, D. J. Horen, and C. Mahaux, Phys. Rev. C **36**, 2252 (1987).

[10] J. Boguta and H. Stöcker, Phys. Lett. **120B**, 289 (1983).

[11] B. Friedman and V. R. Pandharipande, Nucl. Phys. **A361**, 502 (1981).

[12] J. Boguta and S. Bohrmann, Phys. Lett. **102B**, 93 (1981).

[13] M. Rufa, H. Stöcker, P. G. Reinhardt, J. Maruhn, and W. Greiner, J. Phys. G **13**, L143 (1987).

[14] J. H. Taylor and J. M. Weisberg, Princeton University Observatory report, 1989 (to be published).

[15] G. Baym and S. A. Chin, Nucl. Phys. **A262**, 527 (1976). See also J. I. Kapusta, *Finite Temperature Field Theory* (Cambridge Univ. Press, New York, 1989).

[16] J. L. Friedman, J. R. Ipser, and L. Parker, Phys. Rev. Lett. **62**, 3015 (1989).

[17] J. Kristian *et al.*, Nature (London) **338**, 234 (1989).