Equation of state and hybrid star properties with the weakly interacting light U-boson in relativistic models (Postprint)

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Abstract

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Full Text

Preamble

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It has long been a puzzle whether quarks may exist in the interior of massive neutron stars, since the hadron-quark phase transition softens the equation of state (EOS) and reduces the neutron star (NS) maximum mass very significantly. In this work, we consider the light U-boson that increases the NS maximum mass appreciably through its weak coupling to fermions. The inclusion of the U-boson may thus allow the existence of quark degrees of freedom in the interior of large mass neutron stars. Unlike the consequence of the U-boson in hadronic matter, the stiffening role of the U-boson in the hybrid EOS is not sensitive to the choice of the hadron phase models. In addition, we have also investigated the effect of the effective QCD correction on the hybrid EOS. This correction may reduce the coupling strength of the U-boson that is needed to satisfy NS maximum mass constraints.

While the inclusion of the U-boson also increases the NS radius significantly, we find that appropriate in-medium effects of the U-boson may reduce the NS radii significantly, satisfying both the NS radius and mass constraints well.

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INTRODUCTION

The equation of state (EOS) of isospin asymmetric nuclear matter is of prime importance for the investigation of nuclear structure [1–4], heavy-ion reaction dynamics [5, 6], and many issues in astrophysics [6–11]. However, the EOS of asymmetric matter at supra-normal densities is still poorly known [6, 12], though the constraint on the symmetric part of the EOS at supra-normal densities can be extracted from the collective flow data of high energy heavy-ion reactions [13].

Different models and approaches can produce rather different high-density behaviors [14–17], while complexity arises for the EOS of asymmetric matter when phase transitions—such as hyperon production, meson condensation, and quark deconfinement—take place at high densities. Normally, phase transitions can soften the EOS and reduce the neutron star (NS) maximum mass significantly.

Recently, massive NSs have been identified through high-precision measurements, for instance the pulsars PSR J1614-2230 with $M = 1.97 \pm 0.04 M_{\odot}$ [18] and PSR J0348+0432 with $M = 2.01 \pm 0.04 M_{\odot}$ [19]. The larger NS mass implies a stiffer EOS of NS matter at high densities. Considering the phase transitions in the NS core and the EOS constraint from collective flow data at high densities, it is difficult for nuclear models to reproduce NSs as massive as $2M_{\odot}$. This indeed poses a challenge to the nuclear EOS at supra-normal densities. One may doubt whether there are new degrees of freedom besides nucleons in NSs [18, 20, 21]. On the other hand, one may contemplate what interactions can allow these new degrees of freedom in massive NSs [22–35].

The mixture of low-density hadronic matter and high-density quark matter may form in the NS core after the hadron-quark transition occurs with increasing density. In order to evade the potential conflict between the softened EOS due to the appearance of the quark phase and the observed massive NSs, it is necessary to stiffen the quark EOS, for instance by considering strong coupling and/or color superconductivity [24–29, 32]. Recently, a new repulsion provided by the U-boson was introduced for nucleons in NSs [36-38]. The light U-boson, first proposed by Fayet [39], might be regarded as the mediator of the putative fifth force [40–42]. Recently, this light U-boson has been considered as the interaction mediator of MeV dark matter to account for the bright 511 keV γ-ray from the galactic bulge [43–46]. The coupling of the U-boson with nucleons is very weak, but can increase appreciably the NS maximum mass [36–38] and effectively influence the transition density at the inner edge separating the liquid core from the solid crust [47].

In particular, the weak coupling to baryons plays a striking role in stabilizing neutron stars in the presence of super-soft symmetry energy [37] that is extracted from the FOPI/GSI data on the π^-/π^+ ratio in relativistic heavy-ion collisions [48].

In this work, we consider the effect of the U-boson on the softened EOS due to the hadron-quark phase transition. We adopt the relativistic mean-field (RMF) theory, which has achieved great success in the past few decades [49–58], to describe hadronic matter, and the MIT bag model for quark matter [59, 60]. For the hadron-quark phase transition, we use Gibbs construction [61, 62] to determine the mixed phase of hadronic and quark matter. As the stiffening role of the U-boson depends on the softness of the models [38], we will examine how the U-boson stiffens the hybrid EOSs initiated with different hadronic models. Since quarks in the bag model are free of interaction, it is also interesting to investigate briefly the effect of the effective correction from perturbative QCD [26, 63–66]. Eventually, we will investigate the properties of hybrid stars with various EOSs and discuss how the mass and radius constraints from NS observations can be satisfied.

The paper is organized as follows. In Sec. II, we present briefly the formalism. In Sec. III, numerical results and discussions are presented. Finally, a summary is given in Sec. IV.

II. FORMALISM

In the RMF models adopted in this work, the nucleon-nucleon interaction is realized via the exchange of three mesons: the isoscalar meson σ, which provides the intermediate-range attraction between nucleons; the isoscalar-vector meson ω, which offers the short-range repulsion; and the isovector-vector meson ρ, which accounts for the isospin dependence of the nuclear force. Though the π meson interacts strongly with nucleons, we do not include it here because the RMF framework only has the Hartree term that gives zero contribution for pseudoscalar mesons. The π mesons can be included in the relativistic Hartree-

Fock approximation. Without the exchange terms, the RMF approximation still works well due mainly to the fact that in the relativistic framework the fermions are already described by the Dirac equation. Moreover, the interaction in the RMF approximation, built upon the balance between intermediate-range attraction and short-range repulsion, can yield the saturation properties of nuclear matter very well [49-55].

The relativistic Lagrangian is then written as:

$$\mathcal{L} = \bar{\psi}[i\gamma^\mu\partial_\mu - M + g_\sigma\sigma - g_\omega\gamma^\mu\omega_\mu - g_\rho\gamma^\mu\tau_3\rho_\mu]\psi + \mathcal{L}_U + \frac{1}{2}(\partial_\mu\sigma\partial^\mu\sigma - m_\sigma^2\sigma^2) - \frac{1}{4}\omega_{\mu\nu}\omega^{\mu\nu} + \frac{1}{2}m_\omega^2\omega_\mu\omega^\mu - \frac{1}{4}\rho_{\mu\nu}\rho^{\mu\nu} + \frac{1}{2}m_\rho^2$$

where ψ, σ, ω, and ρ are the fields of the nucleon, scalar, vector, and neutral isovector-vector mesons, with masses M, m_σ, m_ω, and m_ρ, respectively. g_i $(i = \sigma, \omega, \rho)$ are the corresponding meson-nucleon couplings. $\omega_{\mu\nu}$ and $\rho_{\mu\nu}$ are the strength tensors of ω and ρ mesons, respectively, with $\omega_{\mu\nu} = \partial_\mu\omega_\nu - \partial_\nu\omega_\mu$ and $\rho_{\mu\nu} = \partial_\mu\rho_\nu - \partial_\nu\rho_\mu$.

The self-interacting terms of σ, ω mesons and the isoscalar-isovector coupling are given generally as:

$$U_{\text{eff}}(\sigma, \omega_\mu, \rho_\mu) = \frac{1}{3}g_2\sigma^3 + \frac{1}{4}g_3\sigma^4 + \frac{1}{4}c_3(\omega_\mu\omega^\mu)^2 + 4\Lambda_V g_\rho^2\omega_\mu\omega^\mu\rho_\mu\rho^\mu$$

In addition, we include in the Lagrangian $\mathcal{L}_U$ for the U-boson that is beyond the standard model. Following the form of the vector meson, $\mathcal{L}_U$ is written as:

$$\mathcal{L}_U = -\bar{\psi}g_u\gamma^\mu u_\mu\psi - \frac{1}{4}U_{\mu\nu}U^{\mu\nu} + \frac{1}{2}m_u^2 u_\mu u^\mu$$

with u_μ the field of the U-boson. $U_{\mu\nu}$ is the strength tensor of the U-boson, $U_{\mu\nu} = \partial_\mu u_\nu - \partial_\nu u_\mu$.

With the standard Euler-Lagrange formalism, we can deduce from the Lagrangian the equations of motion for the nucleon and meson fields. While in the mean-field approximation the Dirac field of nucleons is quantized [52], the fields of mesons and U-boson, which are replaced by their classical expectation values, obey the following equations:

$$m_\sigma^2\sigma = g_\sigma\rho_s - g_2\sigma^2 - g_3\sigma^3,$$
$$m_\omega^2\omega_0 = g_\omega\rho_B - c_3\omega_0^3 - 8\Lambda_V g_\rho^2\omega_0\rho_0^2,$$
$$m_\rho^2\rho_0 = g_\rho\rho_3 - 8\Lambda_V g_\rho^2\omega_0^2\rho_0,$$
$$m_u^2 u_0 = g_u\rho_B,$$

where the temporal subscript of the ρ meson is neglected for convenience, ρ_s and ρ_B are the scalar and baryon densities, respectively, and ρ_3 is the difference between the proton and neutron densities, namely $\rho_3 = \rho_p - \rho_n$. The set of coupled equations can be solved self-consistently using the iteration method. With these mean-field quantities, the resulting energy density ε and pressure P for the hadronic phase are written as:

$$\varepsilon_H = \sum_{i=p,n} \frac{2}{(2\pi)^3} \int d^3k\, E_i^* + \frac{1}{2} m_\sigma^2 \sigma^2 + \frac{1}{3} g_2 \sigma^3 + \frac{1}{4} g_3 \sigma^4 + \frac{1}{2} m_\omega^2 \omega_0^2 + \frac{3}{4} c_3 \omega_0^4 + \frac{1}{2} m_\rho^2 \rho_0^2 + 6\Lambda_V g_\rho^2 \omega_0^2 \rho_0^2 + \frac{1}{2} m_u^2 u_0^2,$$

$$P_H = \sum_{i=p,n} \frac{2}{3(2\pi)^3} \int d^3k \frac{k^2}{E_i^*} - \frac{1}{2} m_\sigma^2 \sigma^2 - \frac{1}{3} g_2 \sigma^3 - \frac{1}{4} g_3 \sigma^4 + \frac{1}{2} m_\omega^2 \omega_0^2 + \frac{1}{4} c_3 \omega_0^4 + \frac{1}{2} m_\rho^2 \rho_0^2 + 2\Lambda_V g_\rho^2 \omega_0^2 \rho_0^2 + \frac{1}{2} m_u^2 u_0^2,$$

with $E_i^* = \sqrt{k^2 + (M_i^*)^2}$.

For the quark phase, we use the MIT bag model, in which the unique parameter —the bag constant—arises from the energy difference between the perturbative ground state and the chiral symmetry breaking vacuum of the theory [67]. In addition to its simplicity, the MIT bag model has been widely used because of its success in describing the vacuum properties of hadrons. As the hadron-quark phase transition occurs, the hadronic degrees of freedom start to turn into quarks which are free of interactions in the bag model. The pressure and energy density are given as:

$$\varepsilon_Q = B + \sum_f \frac{3}{\pi^2} \int_0^{k_F} dk\, k^2 \sqrt{k^2 + m_f^2},$$

$$P_Q = -B + \sum_f \frac{1}{\pi^2} \int_0^{k_F} dk\, k^2 \left(\sqrt{k^2 + m_f^2} - \frac{m_f^2}{\sqrt{k^2 + m_f^2}} \right),$$

where B is the bag constant, and the sum runs over flavor f. With perturbative QCD corrections included to first order [26, 63–66], the grand thermodynamic potentials for quarks are given as:

$$\Omega_u = -\frac{\mu_u^4}{4\pi^2}(1-c),$$

$$\Omega_d = -\frac{\mu_d^4}{4\pi^2}(1-c),$$

$$\Omega_s = \Omega_s(\mu_s, m_s, c),$$

where the term linear in $c = 2\alpha_s/\pi$ is from the perturbative QCD correction, and the masses of up and down quarks are set to zero in obtaining the above expressions. Due to the nonzero strange quark mass ($m_s = 150$ MeV in the calculations of this work), the expression for Ω_s is somewhat tedious and can be referred to Ref. [63]. All thermodynamic quantities follow consistently from Ω. Since quark matter in hybrid stars is not in the perturbative regime, we regard c and the bag constant as effective parameters in the present model, and denote the effective perturbative QCD correction simply as the QCD correction in the following. The pressure is given by $P(\mu) = -\Omega(\mu)$, the quark number density by $n(\mu) = \partial P/\partial\mu$, and the energy density by $\varepsilon = -P + \mu n$.

We use Gibbs construction [61, 62] to describe the hadron-quark phase transition. After the hadron-quark phase transition sets in, one may construct a mixed phase of hadronic and quark matter over a finite range of pressures and densities according to the Gibbs conditions for phase equilibrium. The Gibbs conditions for chemical and mechanical equilibrium and the charge neutrality condition are written as:

$$\begin{gathered}\mu_{b,i}^H - \mu_{c,i}^Q = \mu_i,\\ P_H = P_Q,\\ (1-Y)\rho_b q_b + \sum_f \rho_f q_f + \sum_l \rho_l q_l = 0,\end{gathered}$$

where i runs over baryons, leptons ($b_i = 0$ for leptons), and quarks, and Y is the baryon number fraction of the quark phase. With these conditions, the onset density of the hadron-quark phase transition can be obtained, and the mixed phase is then constructed. Eventually, the total baryon density, energy density, and pressure of the mixed phase are in turn given by:

$$\begin{gathered}\rho_M = (1-Y)\rho_H + Y\rho_Q,\\ \varepsilon_M = (1-Y)\varepsilon_H + Y\varepsilon_Q + \varepsilon_l,\\ P_M = (1-Y)P_H + YP_Q + P_l.\end{gathered}$$

Note that one may also use Maxwell construction for the phase transition. However, Maxwell construction does not produce a mixed phase. In the Maxwell construction, a direct transition from hadronic to quark matter is accompanied by a density jump, and both phases are separately charge neutral. Regardless of charge chemical equilibrium, only a single (baryon) chemical potential is common to the hadronic and quark phases at constant pressure. To avoid charge motion due to different charge chemical potentials in the two phases in the Maxwell construction, a more realistic equation of state can be obtained from the Wigner-Seitz cell calculation by taking into account Coulomb and surface effects. Then, the equation of state of the mixed phase obtained from this approach becomes close to that of the Gibbs construction [34].

Using the EOS of hybrid star matter as input, we may obtain NS properties by solving the Tolman-Oppenheimer-Volkoff (TOV) equation [68, 69]:

$$\frac{dP(r)}{dr} = -\frac{[P(r)+\varepsilon(r)][M(r)+4\pi r^3 P(r)]}{r(r-2M(r))},$$

$$\frac{dM(r)}{dr} = 4\pi r^2 \varepsilon(r),$$

where r is the radial coordinate from the center of the star, $P(r)$ and $\varepsilon(r)$ are the pressure and energy density at position r, respectively, and $M(r)$ is the mass contained within a sphere of radius r. Note that here we use units for which the gravitational constant is $G = c = 1$. The radius R and mass $M(R)$ of a NS are obtained from the condition $P(R) = 0$. Because NS matter undergoes a phase transition from homogeneous to inhomogeneous matter in the low-density region, the EOS obtained from homogeneous matter does not apply to the low-density region, and the empirical low-density EOS from the literature [70, 71] is adopted.

III. RESULTS AND DISCUSSIONS

For the hadronic phase in hybrid stars, we consider the simple composition: neutrons, protons, electrons, and muons. We do not include hyperons in this work. The appearance of hyperons can largely soften the equation of state, thus reducing the mass of neutron stars greatly. Due to accurate mass measurements of large-mass neutron stars, some have even concluded that hyperonic EOSs must be ruled out [18]. Indeed, the onset densities of hyperons are very model-dependent. With the RMF parameter set NL3, the Λ hyperon appears at about 0.28 fm^{-3} [72, 73], while it appears at about 0.48 fm^{-3} with the extended MDI interaction [74]. The superfluidity of hyperons renders the Λ hyperons to appear at around 0.64 fm^{-3} [75], which only affects the properties of neutron stars slightly. When hyperonic degrees of freedom are included, it is found that the hadron-quark transition density cannot change continuously with the emergence of hyperons. This means that the appearance of hyperons disfavors the hadron-quark phase transition in RMF models. A similar effect of hyperons is found in the literature [76]. As implied by large-mass NSs, the in-medium hyperon potential should be density-dependent [23]. We may further explore the density-dependent interaction of hyperons in the future, but this is beyond the scope of the present work.

Among numerous nonlinear RMF models, we select two typical best-fit parameter sets, NL3 [77] and FSUGold [78], to describe the hadron phase of hybrid star matter. Parameters and saturation properties of these two parameter sets are listed in [TABLE:I]. Usually, nonlinear RMF models include nonlinear self-interacting meson terms to simulate appropriate in-medium effects of the strong interaction. The parameter set NL3 includes the nonlinear self-interaction of the σ meson, while FSUGold includes, in addition to the latter, the nonlinear self-interaction of the ω meson. The resulting EOS of FSUGold is much softer than that of NL3 at high densities.

It is known that the vector U-boson can stiffen the EOS, and the stiffening role is much more significant for soft EOSs in pure hadronic matter [38]. It is thus interesting to see whether we can observe a similar phenomenon in hybrid star matter when the hadronic phase is described with these two different RMF models.

[TABLE:I] Parameters and saturation properties for the hadron phase models NL3 and FSUGold. Meson masses, incompressibility κ and symmetry energy are in units of MeV, and the density is in units of fm^{-3}.

While it remains an open question to determine the density at which the hadron-quark phase transition occurs, we investigate the onset of the phase transition in two ways: one is to fix the transition density ρ_c by adjusting the bag constant, and the other is to fix the bag constant. Given the bag constant, we also obtain the transition density which is now dependent on the hadron phase models. Once the transition density is determined, we can construct the mixed phase with the Gibbs conditions and then obtain the EOS of hybrid star matter.

In this work, we choose the bag constant to be $B^{1/4} = 180 \sim 220$ MeV, which is a reasonable range between those values obtained by fitting light-hadron spectra [79–82] and those (e.g., 250 MeV) used in the hydrodynamical evolution of quark-gluon plasma [83, 84]. The present range is also close to that used in the literature [32, 74, 85]. For the choice of transition density, we refer to the literature where it is around $1.5 \sim 4\rho_0$ [26, 86–88]. In addition to the fact that the EOS of hybrid star matter with a transition density of $4\rho_0$ or even higher has only a minor effect on the star' s maximum mass, we choose the range $(2 \sim 3\rho_0)$ for the transition density in this work.

As an example, we calculate here the various phase boundaries with $B^{1/4} = 180$ MeV and with transition densities $2\rho_0$ and $3\rho_0$. For the parameter set NL3, the transition density with $B^{1/4} = 180$ MeV from the hadronic phase to the mixed phase is about 0.20 fm^{-3}, and 0.77 fm^{-3} from the mixed phase to pure quark phase. The extent of the mixed phase is about 0.57 fm^{-3}. For the parameter set FSUGold, the extent of the mixed phase is 1.47 fm^{-3} with onset transition density 0.28 fm^{-3}. For the case with transition density fixed at $2\rho_0$, the extent of the mixed phase is about 0.95 fm^{-3} with parameter set NL3, while it is 1.7 fm^{-3} with FSUGold. For transition density at $3\rho_0$, the extents of the mixed phase are 1.99 fm^{-3} and 2.34 fm^{-3} with NL3 and FSUGold, respectively.

The phase transition usually tends to soften the EOS, consistent with the requirement of spontaneous stability in natural processes. This is also true for the hadron-quark phase transition in isospin asymmetric nuclear matter. In [Figure 1: see original paper], we depict the EOS of hybrid star matter for various cases with onset density $3\rho_0$ for hadron-quark phase transition.

In comparison to the EOS without phase transition as shown in [Figure 1: see original paper], we see that the EOS of isospin asymmetric matter is greatly softened by the hadron-quark phase transition. It is striking to see that the stiff EOS with NL3, which is not favored by constraints from flow data of heavy-ion reactions [13], becomes even much softer than the soft FSUGold EOS due to the hadron-quark transition. With the inclusion of the U-boson, the soft EOS is stiffened greatly. Because the EOS of hybrid star matter initiated with NL3 is now much softer than that with FSUGold, the stiffening effect turns out to be much more appreciable for the EOS initiated with NL3. A similar stiffening role of the U-boson can also be clearly seen in the case without phase transition, as we compare curves with NL3 and FSUGold. Since the softening of the EOS due to the hadron-quark phase transition largely reduces the maximum mass of NSs, we will see below that the inclusion of the U-boson plays an important role in satisfying the maximum mass constraint for NSs.

To see the role of the U-boson specifically, we depict in [Figure 2: see original paper] the EOSs of hybrid star matter for a set of ratio parameters $c_u = (g_u/m_u)^2$. Similar to that shown in [Figure 1: see original paper], the EOS of hybrid star matter is stiffened significantly due to the repulsion provided by the U-boson. We see that the EOSs with the soft (FSUGold) and stiff (NL3) hadron phase models acquire similar stiffening, especially for the case with fixed bag constant. This is different from the case in pure hadronic matter where a much more significant stiffening effect is produced by the soft model [38]. For the case with fixed transition density, the stiffening role of the U-boson is more significant for the stiff parameter set NL3. The reason for this lies in the following facts. In pure hadronic matter with RMF models, there is cancellation between the repulsion provided by the vector meson and the attraction provided by the scalar meson. Thus, more significant cancellation in the soft model sharpens the importance of the repulsion provided by the U-boson. While quarks are modeled by the MIT bag model with the same bag constant, the repulsion provided by the U-boson plays the same stiffening role in the quark phase EOS after the hadron-quark phase transition occurs. In the case with fixed transition density, the quark phase EOS connected to the hadronic phase with NL3 is softer than that with FSUGold because of the larger bag constant. Meanwhile, due to phase equilibrium in the mixed phase, the EOS initiated with NL3 is softer than that with FSUGold with increasing density. Thus, the U-boson provides a more significant stiffening role in the EOS initiated with the stiff NL3 in the hadron phase.

Since in pure hadronic matter the stiffening effect of the U-boson is relevant to the extent of softness of the EOSs, it is interesting to examine whether such dependence exists in the EOS of hybrid star matter with quark degrees of freedom. Shown in [Figure 3: see original paper] are the EOSs of hybrid star matter with and without the contribution of the U-boson for two RMF parameter sets. We can see that after the occurrence of the hadron-quark phase transition, the difference in EOSs with fixed transition density seems more apparent than that with fixed bag constant. This is understandable because different bag constants are used to obtain the same transition density. However, compared to the large difference in the EOSs in pure hadronic matter as shown in [Figure 1: see original paper], we see that the phase transition reduces this difference largely in EOSs at high densities. Compared to the evolution of EOSs with NL3 and FSUGold, we see that the stiffer EOS undergoes more appreciable softening in the mixed phase, as the quark phase is given by the same MIT bag model either with fixed bag constant or with a given transition density. We have also examined results for bag constants ranging from 180 MeV to 220 MeV and transition densities from $3\rho_0$ to $4\rho_0$. The increase of the bag constant or transition density gives rise to a larger extent of the mixed phase, consistent with findings in Ref. [89]. For instance, with the parameter set NL3, the transition density with $B^{1/4} = 200$ MeV is 0.28 fm^{-3}, and the extent of the mixed phase increases from 0.57 fm^{-3} (with $B^{1/4} = 180$ MeV) to 1.1 fm^{-3}. With the parameter set FSUGold, the extent of the mixed phase shows a more apparent rise, reaching 2.8 fm^{-3} starting from transition density 0.47 fm^{-3}. In any case, the appreciable stiffening role of the U-boson in soft EOSs can be clearly observed as in [Figure 3: see original paper], similar to that in [Figure 2: see original paper].

In this work, we make a comparative study with stiff and soft EOSs. It is worth noting that the parameter Λ_V may also affect the EOS of asymmetric matter. This parameter does not appear in the NL3 parameter set. In the FSUGold parameter set, we may change this parameter (along with g_ρ) while keeping the symmetry energy at saturation density unchanged. Here, we choose three parameter sets: FSUGw15 ($\Lambda_V = 0.015$), FSUGold ($\Lambda_V = 0.03$), and FSUGw45 ($\Lambda_V = 0.045$), all of which can fit the ground-state properties of ^{208}Pb [3]. The transition density increases moderately with the rise of Λ_V, since the stiffness of the symmetry energy controlled by Λ_V can affect the chemical equilibrium. With bag constant $B^{1/4} = 180$ MeV, we obtain the following transition densities: 0.252, 0.282, and 0.297 fm^{-3} with parameter sets FSUGw15, FSUGold, and FSUGw45, respectively.

In the above, the quarks with the bag model are free of interactions. As the QCD coupling is not negligible at densities of interest for compact star physics, it is necessary to discuss briefly the QCD correction. It is reasonable to take the value of c in the equations above to be 0.3 [26]. Shown in [Figure 4: see original paper] is the EOS with the QCD correction. We see that the QCD correction can moderately stiffen the EOS of hybrid star matter at high densities. The QCD correction may affect the phase transition and increase the extent of the mixed phase due to the fact that the quark phase pressure is now modified by the QCD correction. For instance, with bag constant $B^{1/4} = 180$ MeV, the onset density of the mixed phase with parameter set NL3 is about 0.28 fm^{-3} with the extent of the mixed phase being about 0.96 fm^{-3}. With parameter set FSUGold, the onset density of the mixed phase is about 0.46 fm^{-3}, and the extent of the mixed phase grows dramatically up to 6.49 fm^{-3}. As observed in [Figure 4: see original paper], the QCD correction to the EOS is moderately dependent on the hadron phase models. Specifically, the QCD correction results in an extent of the mixed phase with the soft FSUGold being larger than that with the stiff NL3. We should note that this specific dependence of the QCD correction on the hadron phase models is due to the connection built upon the phase equilibrium conditions. For comparison, we also plot results with and without the U-boson. As shown in [Figure 4: see original paper], we see that the QCD correction is just moderate compared with the contribution of the U-boson. The stiffening role of the U-boson is similar in cases with and without the QCD correction.

Now we turn to investigations of hybrid star properties with these EOSs. Shown in [Figure 5: see original paper] is the relationship between NS mass and radius for various EOSs with the inclusion of the U-boson. For the stiff hadron phase EOS, the hadron-quark transition at high density does not reduce the maximum mass of hybrid stars significantly, as shown in the lower right panel of [Figure 5: see original paper]. In most cases, however, the phase transition may result in very significant reduction of the maximum mass, which would be inconsistent with observations of large-mass NSs [18, 19]. We see that the inclusion of the U-boson can remedy this mass deficit very efficiently with an appropriately chosen ratio parameter. The rise of the maximum mass is consistent with the corresponding stiffening of the high-density EOS shown in [Figure 2: see original paper]. We can see that the role of the U-boson in increasing the maximum mass is more significant for softer EOSs.

Though only results with particular bag constant and transition density are shown in [Figure 5: see original paper], it is sufficient to reveal the role of the U-boson since changing these parameters yields rather similar results. It is, however, interesting to examine the consequences for the internal structure of particular NSs by varying these parameters. Let us first compare results with $B^{1/4} = 180$ and 200 MeV. With $B^{1/4} = 180$ MeV, the parameter set NL3 gives a NS maximum mass of $1.39M_\odot$. The size of the quark core is 5.76 km, the mixed phase spreads from 5.76 to 7.62 km, and the hadronic phase extends from 7.62 to 9.68 km. With $B^{1/4} = 200$ MeV, the maximum mass is $1.86M_\odot$. For convenient comparison, we also examine the internal structure of the $1.39M_\odot$ NS with $B^{1/4} = 200$ MeV. The size of the quark core is now 4.42 km, the mixed phase spreads from 4.42 to 6.38 km, and the hadronic phase extends from 6.38 to 10.59 km. This indicates that the smaller bag constant, which gives rise to a smaller transition density, features a larger quark core in a particular NS. For the case with smaller transition density, the conclusion is similar since the smaller transition density is given by the smaller bag constant.

It is also interesting to examine the effect of the U-boson on NS structure. Although the inclusion of the U-boson does not change the chemical and mechanical equilibria, the transition densities, or the extents of each phase, it does affect the relative proportions. As an example, we examine the size of the quark core and mixed phase for the maximum mass configuration with parameters $B^{1/4} = 180$ MeV and $(g_u/m_u)^2 = 25$ GeV^{-2}. With the inclusion of the U-boson, the NS maximum mass increases from $1.39M_\odot$ to $1.71M_\odot$. In the NS with this maximum mass configuration, the radius of the pure quark core is 5.3 km, which is smaller than 5.76 km obtained without the U-boson. The shell of the mixed phase extends from 5.3 to 8.6 km, while without the U-boson, the shell extends from 5.76 to 7.6 km. These results indicate that the U-boson may change the ratio of various phases in a particular NS, though it does not directly participate in the phase transition.

As shown in [Figure 5: see original paper], the NS radius is significantly increased by the U-boson. It was pointed out in the literature [6, 7] that the NS radius is primarily determined by the EOS in the lower density region of $1\rho_0$ to $2\rho_0$. Since the inclusion of the U-boson also increases the pressure in the lower density region appreciably, a large NS radius is obtained accordingly with various c_u values of interest. This is similar to previous works in Refs. [37, 38]. It is known that extraction of NS radii from observations still suffers from large systematic uncertainties [90] involved in distance measurements and theoretical analyses of the light spectrum [8, 91–93]. There is yet no consensus on extracted NS radii to date. For instance, using thermal spectra from quiescent low-mass X-ray binaries (qLMXBs), Guillot and collaborators extracted NS radii of $R_{NS} = 9.4 \pm 1.2$ km [94], while a relevant study of spectroscopic radius measurements also suggests small radii of $10.8^{+0.5}_{-0.4}$ km for a $1.5M_\odot$ NS [95]. There are also larger extracted radii: a 3σ lower limit of 11.1 km on the radius of PSR J0437-4715 [96] and a lower limit of 13 km for 4U 1608-52 [97]. If small NS radii are established, we need to reconsider the large NS radii produced by the U-boson. A possible solution for NS radius suppression is to invoke appropriate in-medium effects [98]. Considering that NS radii are mainly determined by the low-density EOS, it is possible to reduce NS radii by constructing a density-dependent coupling constant for the U-boson in the low-density regime.

In the high-density regime, we neglect the in-medium effect for the U-boson, as one knows that the in-medium effect at high densities is greatly suppressed by Pauli blocking [98, 100, 101]. As an example, we perform calculations with the EOS whose hadron phase part is obtained with FSUGold and find that an appropriate density-dependent coupling constant of the U-boson can reduce NS radii significantly, as shown in [Figure 6: see original paper]. The density-dependent coupling constant, depicted in the inset of [Figure 6: see original paper], is designed to change the pressure little in the low-density regime [98]. While the energy density still acquires a significant increase from the U-boson, the resulting NS radius is even less than that obtained with the pure hadron phase EOS. Since the NS maximum mass is determined mainly by the high-density EOS, the present form of density dependence just gives rise to a moderate reduction of the NS maximum mass.

In [Figure 5: see original paper] and [Figure 6: see original paper], we have not included the contribution of the QCD correction. The influence of the QCD correction on the NS maximum mass is rather sensitive to the transition density. Given large transition densities, the QCD correction plays only a limited role in enhancing the NS maximum mass, since the maximum mass in this case is dominated by the hadron phase EOS. As the transition density decreases, the QCD correction brings more significant enhancement to the NS maximum mass. These findings are consistent with those in Ref. [32]. On the other hand, when we fix the bag constant, the situation for the QCD correction is different because the transition density itself is increased by the correction through the phase equilibrium conditions. As shown in [Figure 7: see original paper], the QCD correction with fixed bag constant gives rise to significant enhancement of the NS maximum mass, which is mainly attributed to the appreciable rise of the transition density. In this case, since the hadron phase EOS with NL3 is much stiffer than that with FSUGold, the large transition density due to the QCD correction leads to more appreciable enhancement of the NS maximum mass with NL3. With the inclusion of the U-boson, the NS maximum mass can further increase to satisfy the $2M_{\odot}$ constraint [18, 19]. It is interesting to see that we need only a smaller ratio parameter of the U-boson to meet the maximum mass constraint in this case. This is more apparent for the EOS with the stiff NL3: the ratio parameter is largely reduced by the QCD correction, as shown in [Figure 7: see original paper].

Since the magnitude and in-medium behavior of the U-boson ratio parameter are important for NS properties, it remains significant to discuss the parameters for the U-boson. To satisfy the NS maximum mass constraint, the ratio parameter c_u in this work is estimated to be around $0 \sim 50$ GeV^{-2}, depending on the hadron phase models chosen. To avoid violation of low-energy nuclear constraints for finite nuclei, we may limit the weak interaction strength of the U-boson with mass being of order 1 MeV [38, 98]. For instance, if g_u is 0.01, then the mass of the U-boson would be below 1.4 MeV, responsible for a long-range interaction. The weak interaction strength of the U-boson does not compromise the success of nuclear models in reproducing properties of finite nuclei [99]. Interestingly, the inclusion of the QCD correction may reduce the ratio parameter significantly. We note that these estimated orders of magnitude for the U-boson parameters can be compatible with parameter regions allowed by a few experimental constraints [36]. Moreover, the density-dependent parameter of the U-boson is found to be consistent with usual predictions on the NS radius, while the density dependence would originate from in-medium effects in the nuclear many-body system [98, 100, 101].

IV. SUMMARY

In this work, we have investigated the hadron-quark phase transition using Gibbs conditions with RMF models for the hadron phase and the MIT bag model for the quark phase. We have considered the U-boson to stiffen the EOS of hybrid matter that is greatly softened by the phase transition.

With the inclusion of the U-boson, the hybrid EOS is appreciably stiffened, and the stiffening extent is similar in EOSs with stiff and soft hadron phase models due to the fact that the same MIT model is used for the quark phase. As a result, the NS maximum mass is significantly increased by the U-boson. In addition, we have investigated the effect of the effective QCD correction on the hybrid EOS. This correction may give rise to stiffening of the hybrid EOS and increase of the NS maximum mass that is significant when the transition density is not very high. The effective QCD correction can reduce the coupling strength of the U-boson needed to satisfy NS maximum mass constraints. While the inclusion of the U-boson also increases the NS radius significantly, we find that appropriate density dependence of the U-boson coupling constant may bring the NS radius down to a regime consistent with observations and other model predictions. We have also discussed that with weak interaction strength, the inclusion of the U-boson does not compromise the success of conventional nuclear models in reproducing properties of finite nuclei. In summary, the inclusion of the U-boson can favorably allow the existence of quark degrees of freedom in the NS interior that would otherwise be declined by large-mass NS observations. On the other hand, future coincident measurements and more precise extraction of the mass and radius of neutron stars may also constrain the U-boson that is yet to be detected in terrestrial laboratories.

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