Properties of hyperonic matter in strong magnetic fields

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Abstract

We study the effects of strong magnetic fields on the properties of hyperonic matter. We employ the relativistic mean field theory, which is known to provide excellent descriptions of nuclear matter and finite nuclei. The two additional hidden-strangeness mesons, $\sigma^*$ and $\phi$, are taken into account, and some reasonable hyperon potentials are used to constrain the meson-hyperon couplings, which reflect the recent developments in hypernuclear physics. It is found that the effects of strong magnetic fields become significant only for magnetic field strength $B > 5 \times 10^{18}$ G. The threshold densities of hyperons can be significantly altered by strong magnetic fields. The presence of hyperons makes the equation of state (EOS) softer than that in the case without hyperons, and the softening of the EOS becomes less pronounced with increasing magnetic field strength.

PACS numbers: 26.60.-c, 24.10.Jv, 98.35.Eg

Keywords: Relativistic mean-field theory, Hyperonic matter, Magnetic fields
I. INTRODUCTION

Neutron stars are fascinating celestial objects as natural laboratories for dense matter physics. They are associated with the most exotic environments and violent phenomena of the universe [1]. Observations of ordinary radio pulsars indicate that they possess surface magnetic fields of the order of $10^{12}$ G [2]. Recent surveys of soft-gamma repeaters (SGRs) and anomalous x-ray pulsars (AXPs) imply that the surface magnetic field of young neutron stars could be of order $10^{14}$–$10^{15}$ G [3]. The magnetic field strength may vary significantly from the surface to the center in neutron stars. So far, there is no direct observational evidence for the internal magnetic fields of the star, while it may reach $10^{18}$ G, as estimated in some theoretical works [4, 5, 6]. There are many studies of the effects of strong magnetic fields on the properties of dense nuclear matter and neutron stars [2, 4, 5, 6, 7, 8, 9], and the inclusion of hyperons and boson condensation has also been investigated in the literature [10, 11, 12, 13, 14]. On the other hand, much work has been done to investigate the effects of the magnetic field on color superconducting phases of dense quark matter, which are conjectured to exist inside neutron stars [15, 16, 17, 18].

Our knowledge of neutron star interiors is still uncertain. The central density of neutron stars can be extremely high, and many possibilities for such dense matter have been suggested [1, 19]. For densities below twice normal nuclear matter density ($\rho_0 \sim 0.15$ fm$^{-3}$), the matter consists of only nucleons and leptons. When the density is higher than 2$\rho_0$, the equation of state (EOS) and composition of matter are much less certain. The presence of hyperons in neutron stars has been studied by many authors [20, 21, 22, 23, 24, 25]. With increasing density, hyperons may appear through the weak interaction due to the fast rise of the baryon chemical potentials. The presence of hyperons tends to soften the EOS at high density and lower the maximum mass of neutron stars [19, 20, 21, 22, 23, 24, 25]. K$^-$ condensation in dense matter was suggested by Kaplan and Nelson [26] and has been discussed in many works [13, 14, 19, 20]. In general, the chemical potential of antikaons decreases with increasing density because of the interaction with baryons. As a consequence, the ground state of matter at high density may contain antikaon condensation which can soften the EOS of neutron star matter. It has been suggested that the quark matter may exist in the core of massive neutron stars, and the hadron-quark phase transition can proceed through a mixed phase of hadronic and quark matter [1, 19, 22, 27]. If deconfined quark matter does
exist inside stars, it is likely to be in a color superconducting phase \[1, 22\]. At present, the
densities at which these phases occur are still rather uncertain.

In this article, we focus on the properties of hyperonic matter, which is composed of
a chemically equilibrated and charge-neutral mixture of nucleons, hyperons, and leptons.
It is well known that hyperons may appear around twice normal nuclear matter density
and can soften the EOS at high density in the field-free case \[19, 20, 21, 22, 23\]. The
meson-hyperon couplings play an important role in determining the EOS and composition
of hyperonic matter \[24, 25\]. In the presence of strong magnetic fields, the pressure and
composition of matter can be affected significantly. In Ref. \[10\], the authors investigated the
effects of strong magnetic fields on the properties of neutron star matter including hyperons.
They employed two classes of the relativistic mean field (RMF) models, GM1-3 and ZM
models, while the meson-hyperon couplings were assumed to be equal for all hyperons in
their study. It was found that the EOS at high density could be significantly affected both
by Landau quantization and by magnetic moment interactions, but only for field strength
\[B > 5 \times 10^{18} \text{ G} \]. It is very interesting to investigate the influence of strong magnetic fields on
hyperonic matter with recent developments in hypernuclear physics. A recent observation of
the double-Λ hypernucleus \[^{6}_{\Lambda\Lambda}\text{He}\], called the Nagara event \[28\], has had a significant impact
on strangeness nuclear physics. The Nagara event provides unambiguous identification of
\[^{6}_{\Lambda\Lambda}\text{He}\] production with a precise ΛΛ binding energy value \[B_{\Lambda\Lambda} = 7.25 \pm 0.19^{+0.18}_{-0.11} \text{ MeV}\], which
suggests that the effective ΛΛ interaction should be considerably weaker \(\Delta B_{\Lambda\Lambda} \approx 1 \text{ MeV}\)
than that deduced from the earlier measurement \(\Delta B_{\Lambda\Lambda} \approx 4-5 \text{ MeV}\). The weak hyperon-
hyperon \((YY)\) interaction suggested by the Nagara event has been used to reinvestigate the
properties of multistrange systems, and it has been found that the change of YY interactions
affects the properties of strange hadronic matter dramatically \[29, 30, 31\]. We would like to
examine how the YY interaction affects the properties of hyperonic matter in the presence
of strong magnetic fields.

We employ the RMF theory to study the properties of hyperonic matter in the presence
of strong magnetic fields. The RMF theory has been successfully and widely used for the
description of nuclear matter and finite nuclei \[32, 33, 34, 35, 36\]. It has also been applied
to providing the EOS of dense matter for use in supernovae and neutron stars \[37\]. In
the RMF approach, baryons interact through the exchange of scalar and vector mesons.
The meson-nucleon coupling constants are generally determined by fitting to some nuclear
matter properties or ground-state properties of finite nuclei. However, there are large uncertainties in the meson-hyperon couplings due to limited experimental data in hypernuclear physics. Generally, one can use the coupling constants derived from the quark model or the values constrained by reasonable hyperon potentials. The two additional hidden-strangeness mesons, $\sigma^*$ and $\phi$, were originally introduced to obtain the strong attractive $YY$ interaction deduced from the earlier measurement [38]. Their couplings to hyperons are related to the $YY$ interaction [27]. In this study, we adopt the TM1 version of the RMF theory, which includes nonlinear terms for both $\sigma$ and $\omega$ mesons [39]. The TM1 model is known to provide excellent descriptions of the ground states of finite nuclei including unstable nuclei and also shown to reproduce satisfactory agreement with experimental data in the studies of nuclei with deformed configuration and the giant resonances within the random-phase approximation (RPA) formalism [40, 41, 42]. As for the meson-hyperon couplings, we use the values constrained by reasonable hyperon potentials that include the updated information from recent developments in hypernuclear physics. We take into account the two additional hidden-strangeness mesons, $\sigma^*$ and $\phi$, and consider two cases of the $YY$ interaction, the weak and strong $YY$ interactions. By comparing the results with different $YY$ interactions, we evaluate how sensitive the results are to this factor. Because the properties of hyperonic matter are very sensitive to the meson-hyperon couplings, an investigation of the influence of strong magnetic fields on hyperonic matter with a suitable choice of the meson-hyperon couplings would be of interest for the study of neutron stars.

This article is arranged as follows. In Sec. II we briefly describe the RMF theory used to describe hyperonic matter in the presence of strong magnetic fields. In Sec. III, we discuss the parameters used in the calculation. We show and discuss the numerical results in Sec. IV. Section V is devoted to a summary.

II. FORMALISM

We briefly explain the RMF theory used to describe hyperonic matter in the presence of strong magnetic fields. In the RMF approach, baryons interact through the exchange of scalar and vector mesons. The baryons considered in this work are nucleons ($p$ and $n$) and hyperons ($\Lambda$, $\Sigma$, and $\Xi$). The exchanged mesons include isoscalar scalar and vector mesons ($\sigma$ and $\omega$), isovector vector meson ($\rho$), and two additional hidden-strangeness mesons ($\sigma^*$
and $\phi$). In the presence of strong magnetic fields, the total Lagrangian density of matter consisting of a neutral mixture of baryons and leptons takes the form

$$L_{\text{RMF}} = \sum_b \bar{\psi}_b \left[ i\gamma_\mu \partial^\mu - m_b - g_{\sigma b} \sigma - g_{\sigma b}^* \sigma^* 
- g_{\omega b} \gamma_\mu \omega^\mu - g_{\phi b} \gamma_\mu \phi^\mu - g_{\rho b} \gamma_\mu \tau_b \rho_i^\mu 
- q_b \gamma_\mu A^\mu - \frac{1}{2} \kappa_b \sigma_{\mu\nu} F^{\mu\nu} \right] \psi_b 
+ \sum_l \bar{\psi}_l \left[ i\gamma_\mu \partial^\mu - m_l - q_l \gamma_\mu A^\mu \right] \psi_l 
+ \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - \frac{1}{2} m_\sigma^2 \sigma^2 - \frac{1}{3} g_2 \sigma^3 - \frac{1}{4} g_3 \sigma^4 
- \frac{1}{4} W_{\mu\nu} W^{\mu\nu} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu + \frac{1}{4} c_3 (\omega_\mu \omega^\mu)^2 
- \frac{1}{4} R_{\mu\nu} R^{\mu\nu} + \frac{1}{2} m_\rho^2 \rho_\mu \rho_\mu^\mu + \frac{1}{2} \partial_\mu \sigma^* \partial^\mu \sigma^* 
- \frac{1}{4} m_\phi^2 \phi^* \phi^\mu - \frac{1}{2} m_\phi \phi_\mu \phi^\mu - \frac{1}{4} F_{\mu\nu} F^{\mu\nu},$$

(1)

where $\psi_b$ and $\psi_l$ are the baryon and lepton fields, respectively. The index $b$ runs over the baryon octet ($p$, $n$, $\Lambda$, $\Sigma^+$, $\Sigma^0$, $\Sigma^-$, $\Xi^0$, $\Xi^-$), and the sum on $l$ is over electrons and muons ($e^-$ and $\mu^-$). The field tensors of the vector mesons, $\omega$, $\rho$, and $\phi$, are denoted by $W_{\mu\nu}$, $R_{\mu\nu}$, and $S_{\mu\nu}$, respectively. $A^\mu = (0, 0, Bx, 0)$ refers to a constant external magnetic field along the z axis with its field tensor denoted by $F_{\mu\nu}$. It has been found that the contributions from the anomalous magnetic moments of baryons should be considered in strong magnetic fields [4, 5, 6, 8, 9]. The anomalous magnetic moments are introduced via the minimal coupling of the baryons to the electromagnetic field tensor in the term $-\frac{1}{2} \kappa_b \sigma_{\mu\nu} F^{\mu\nu}$, where $\kappa_b = (\mu_b / \mu_N - q_b m_p / m_b) \mu_N$ is the anomalous magnetic moment of baryon species $b$ ($\mu_N = e\hbar / 2m_p$ is the nuclear magneton). As for leptons, the effects of anomalous magnetic moments on the EOS are very small, as shown in Ref. [8], and the electron anomalous magnetic moments could be efficiently reduced by high-order contributions from the vacuum polarization in strong magnetic fields [43]. Therefore, we neglect the anomalous magnetic moments of leptons in the present work.

In the RMF approach, the meson fields are treated as classical fields, and the field operators are replaced by their expectation values. The meson field equations in uniform matter
have the following form:

\[ m_\sigma^2 \sigma + g_2 \sigma^2 + g_3 \sigma^3 = - \sum_b g_{\sigma b} \rho_b^b, \]  
(2)

\[ m_\omega^2 \omega + c_3 \omega^3 = \sum_b g_{\omega b} \beta_b^b, \]  
(3)

\[ m_\rho^2 \rho = \sum_b g_{\rho b} \tau_{3b} \rho_b^b, \]  
(4)

\[ m_\sigma^* \sigma^* = - \sum_b g_{\sigma^* b} \beta_s^b, \]  
(5)

\[ m_\phi^2 \phi = \sum_b g_{\phi b} \rho_b^b, \]  
(6)

where \( \sigma = \langle \sigma \rangle, \omega = \langle \omega^0 \rangle, \rho = \langle \rho_3^0 \rangle, \sigma^* = \langle \sigma^* \rangle, \) and \( \phi = \langle \phi^0 \rangle \) are the nonvanishing expectation values of meson fields in uniform matter. In the presence of strong magnetic fields, the Dirac equations for baryons and leptons are, respectively, given by

\[
\left( i \gamma^\mu \partial_\mu - m_b^* - g_{\omega b} \gamma_0 \omega - g_{\phi b} \gamma_0 \phi - g_{\rho b} \gamma_0 \tau_{3b} \rho \right. \\
\left. -q_b \gamma_\mu A^\mu - \frac{1}{2} \kappa_b \sigma_{\mu \nu} F^{\mu \nu} \right) \psi_b = 0,
\]  
(7)

\[
\left( i \gamma^\mu \partial_\mu - m_l - q_l \gamma_\mu A^\mu \right. \\
\left. -q_b \gamma_\mu A^\mu \right) \psi_l = 0,
\]  
(8)

where \( m_b^* = m_b + g_{\omega b} \sigma + g_{\sigma^* b} \sigma^* \) is the effective mass of baryon species \( b \). Following the method of Refs. \[4, 6, 44\], we solve the Dirac equations and obtain the energy spectra for charged and uncharged baryons as

\[ E_{\nu, s}^b = \sqrt{k_z^2 + \left( \sqrt{m_b^* + 2\nu q_b B - s \kappa_b B} \right)^2 + g_{\omega b} \omega + g_{\phi b} \phi + g_{\rho b} \tau_{3b} \rho}, \]  
(9)

\[ E_s^b = \sqrt{k_z^2 + \left( \sqrt{m_b^* + 2\nu q_b B - s \kappa_b B} \right)^2 + g_{\omega b} \omega + g_{\phi b} \phi + g_{\rho b} \tau_{3b} \rho}, \]  
(10)

while the lepton energy spectrum is expressed by

\[ E_{\nu, s}^l = \sqrt{k_z^2 + m_l^2 + 2\nu q_l B}. \]  
(11)

Here \( \nu = 0, 1, 2, \ldots \) enumerates the Landau levels of charged fermions. The quantum number \( s \) is equal to \( \pm 1 \). For charged baryons, the scalar and vector densities are, respectively, given by \[4, 6\]

\[ \rho_s^b = \frac{|q_b| B m_b^*}{2\pi^2} \sum_{\nu} \sum_s \frac{\sqrt{m_b^* + 2\nu q_b B - s \kappa_b B}}{\sqrt{m_b^* + 2\nu q_b B}} \ln \left| \frac{k_{f, \nu, s}^b + E_{f}^b}{\sqrt{m_b^* + 2\nu q_b B - s \kappa_b B}} \right|, \]  
(12)

\[ \rho_v^b = \frac{|q_b| B}{2\pi^2} \sum_{\nu} \sum_s k_{f, \nu, s}^b, \]  
(13)
where \( k_{f,\nu,s}^b \) is the Fermi momentum of charged baryon \( b \) with quantum numbers \( \nu \) and \( s \). The Fermi energy \( E_f^b \) is related to the Fermi momentum \( k_{f,\nu,s}^b \) by

\[
E_f^{b2} = (k_{f,\nu,s}^b)^2 + \left( \sqrt{m_b^* + 2\nu |q_b| B - s\kappa B} \right)^2.
\]

For uncharged baryons, there is no Landau level quantum number \( \nu \), so the Fermi momentum is denoted by \( k_{f,s}^b \). The Fermi energy \( E_f^b \) is related to the Fermi momentum \( k_{f,s}^b \) by

\[
E_f^{b2} = (k_{f,s}^b)^2 + (m_b^* - s\kappa B)^2.
\]

The scalar and vector densities of uncharged baryon \( b \) are, respectively, given by

\[
\rho_s^b = \frac{m_b^*}{4\pi^2} \sum_s \left[ k_{f,s}^b E_f^b - (m_b^* - s\kappa B)^2 \ln \left( \frac{k_{f,s}^b + E_f^b}{m_b^* - s\kappa B} \right) \right],
\]

\[
\rho_v^b = \frac{1}{2\pi^2} \sum_s \left\{ \frac{1}{3} (k_{f,s}^b)^3 - \frac{1}{2} s\kappa B \left( (m_b^* - s\kappa B) k_{f,s}^b + E_f^b \right) \arcsin \left( \frac{m_b^* - s\kappa B}{E_f^b} - \frac{\pi}{2} \right) \right\}.
\]

For hyperonic matter consisting of a neutral mixture of baryons and leptons, the \( \beta \)-equilibrium conditions without trapped neutrinos are given by

\[
\mu_p = \mu_{\Sigma^+} = \mu_n - \mu_e,
\]

\[
\mu_{\Lambda} = \mu_{\Sigma^0} = \mu_{\Xi^0} = \mu_n,
\]

\[
\mu_{\Sigma^-} = \mu_{\Xi^-} = \mu_n + \mu_e,
\]

\[
\mu_{\mu} = \mu_e,
\]

where \( \mu_i \) is the chemical potential of species \( i \). The chemical potentials of baryons and leptons are, respectively, given by

\[
\mu_b = E_f^b + g_{\omega\omega} \omega + g_{\phi\phi} \phi + g_{\rho\rho} \tau_3 \rho_b,
\]

\[
\mu_l = E_f^l = \sqrt{(k_{f,\nu,s}^l)^2 + m_l^2 + 2\nu |q_l| B}.
\]

The electric charge neutrality condition is expressed by

\[
\sum_b q_b \rho_v^b + \sum_l q_l \rho_v^l = 0,
\]
where the vector density of leptons has a similar expression to that of charged baryons

$$\rho^l_v = \frac{|q_l| B}{2\pi^2} \sum_{\nu, s} k^l_{f, \nu, s}. \quad (25)$$

We solve the coupled Eqs. (2)–(6), (18)–(21), and (24) self-consistently at a given baryon density $\rho_b = \sum_b \rho^b_v$ in the presence of strong magnetic fields. The total energy density of matter is given by

$$\varepsilon = \sum_b \varepsilon_b + \sum_l \varepsilon_l + \frac{1}{2} m^2 \sigma^2 + \frac{1}{3} g_2 \sigma^3 + \frac{1}{4} g_3 \sigma^4 + \frac{1}{2} m^2 \omega^2$$

$$+ \frac{3}{4} c_3 \omega^4 + \frac{1}{2} m^2 \rho^2 + \frac{1}{2} m^2 \sigma^2 + \frac{1}{2} m^2 \phi^2, \quad (26)$$

where the energy densities of charged baryons and leptons have the following forms:

$$\varepsilon_b = \frac{|q_b| B}{4\pi^2} \sum_{\nu, s} \left[ k^{b}_{f, \nu, s} E^{b}_f + \left( \sqrt{m^2_b + 2\nu |q_b| B - s \kappa_b B} \right)^2 \right]$$

$$\times \ln \left[ \frac{k^{b}_{f, \nu, s} + E^{b}_f}{\sqrt{m^2_b + 2\nu |q_b| B - s \kappa_b B}} \right], \quad (27)$$

$$\varepsilon_l = \frac{|q_l| B}{4\pi^2} \sum_{\nu, s} \left[ k^{l}_{f, \nu, s} E^{l}_f + \left( m^2_l + 2\nu |q_l| B \right) \ln \frac{k^{l}_{f, \nu, s} + E^{l}_f}{\sqrt{m^2_l + 2\nu |q_l| B}} \right], \quad (28)$$

while those of uncharged baryons are given by

$$\varepsilon_b = \frac{1}{4\pi^2} \sum_{s} \left\{ \frac{1}{2} k^{b}_{f, s} E^{b}_f - \frac{2}{3} s \kappa_b B E^{b}_f \left( \arcsin \frac{m^*_b - s \kappa_b B}{E^{b}_f} - \frac{\pi}{2} \right) - \left( \frac{s \kappa_b B}{3} + \frac{m^*_b - s \kappa_b B}{4} \right) \right\}$$

$$\times \left( m^*_b - s \kappa_b B \right) k^{b}_{f, s} E^{b}_f + \left( m^*_b - s \kappa_b B \right)^3 \ln \left[ \frac{k^{b}_{f, s} + E^{b}_f}{m^*_b - s \kappa_b B} \right] \right\}. \quad (29)$$

The pressure of the system can be obtained by

$$P = \sum_b \mu_b \rho^b_v + \sum_l \mu_l \rho^l_v - \varepsilon = \mu_n \rho_b - \varepsilon, \quad (30)$$

where the electric charge neutrality and $\beta$-equilibrium conditions are used to get the last equality. We note that the contribution from electromagnetic fields to the energy density and pressure, $\varepsilon_f = P_f = B^2/8\pi$, is not taken into account in the present calculation. In general, the strong magnetic fields in neutron stars can produce magnetic forces that play an important role in determining the structure of the star.
III. PARAMETERS

In this section, we discuss the parameters used in the RMF approach. For the nucleonic sector, we employ the TM1 parameter set which was determined by a least-squares fit to experimental results including stable and unstable nuclei \[39\]. It has been shown that the RMF theory with the TM1 parameter set reproduces satisfactory agreement with experimental data in the studies of the nuclei with deformed configuration and the giant resonances within the RPA formalism \[40, 41, 42\]. With the TM1 parameter set, the nuclear matter saturation density is \(0.145 \text{ fm}^{-3}\), the energy per nucleon is \(-16.3 \text{ MeV}\), the symmetry energy is 36.9 MeV, and the incompressibility is 281 MeV \[39\]. The TM1 parameter set can be found in Refs. \[27, 39\].

The baryons considered in this work are the baryon octet (\(p, n, \Lambda, \Sigma^+, \Sigma^0, \Sigma^-, \Xi^0, \Xi^-\)). In Table I, we list the static properties of baryons such as masses, charges, and magnetic moments. The exchanged mesons include isoscalar scalar and vector mesons (\(\sigma\) and \(\omega\)), isovector vector meson (\(\rho\)), and two additional hidden-strangeness mesons (\(\sigma^*\) and \(\phi\)). As for the meson-hyperon couplings, we take the naive quark model values for the vector coupling constants,

\[
\frac{1}{3}g_{\omega N} = \frac{1}{2}g_{\omega \Lambda} = \frac{1}{2}g_{\omega \Sigma} = g_{\omega \Xi},
\]

\[
g_{\rho N} = \frac{1}{2}g_{\rho \Sigma} = g_{\rho \Xi}, \quad g_{\rho \Lambda} = 0,
\]

\[
2g_{\phi \Lambda} = 2g_{\phi \Sigma} = g_{\phi \Xi} = -\frac{2\sqrt{2}}{3}g_{\omega N}, \quad g_{\phi N} = 0.
\]

The scalar coupling constants are chosen to give reasonable hyperon potentials. We denote the potential depth of the hyperon species \(i\) in the matter of the baryon species \(j\) by \(U_{ij}^{(N)}\). It is estimated from the experimental data of single-\(\Lambda\) hypernuclei that the potential depth of a \(\Lambda\) in saturated nuclear matter should be around \(U_{\Lambda}^{(N)} \simeq -30 \text{ MeV}\) \[43\]. For \(\Sigma\) hyperons, the analysis of \(\Sigma\) atomic experimental data suggests that \(\Sigma\)-nucleus potentials have a repulsion inside the nuclear surface and an attraction outside the nucleus with a sizable absorption. In recent theoretical works, the \(\Sigma\) potential in saturated nuclear matter is considered to be repulsive with the strength of about 30 MeV \[45, 46\]. Recent developments in hypernuclear physics suggest that \(\Xi\) hyperons in saturated nuclear matter have an attraction of around 15 MeV \[46, 47\]. In this article, we use \(U_{\Lambda}^{(N)} = -30\), \(U_{\Sigma}^{(N)} = +30\), and \(U_{\Xi}^{(N)} = -15 \text{ MeV}\) to determine the scalar coupling constants \(g_{\sigma \Lambda}, g_{\sigma \Sigma},\) and \(g_{\sigma \Xi}\), respectively.
The hyperon couplings to the hidden-strangeness meson $\sigma^*$ are restricted by the relation $U^{(\Xi)}_\Xi \simeq 2U^{(\Lambda)}_\Lambda \simeq 2U^{(\Lambda)}_\Lambda$ obtained in Ref. [48]. We consider two cases of hyperon-hyperon ($YY$) interactions. The weak $YY$ interaction implied by the Nagara event suggests $U^{(\Lambda)}_\Lambda \simeq -5$ MeV, while the strong $YY$ interaction deduced from the earlier measurement suggests $U^{(\Lambda)}_\Lambda \simeq -20$ MeV [27, 29, 30, 31]. In Table III we present the meson-hyperon couplings determined by these hyperon potentials. We note that $m_{\sigma^*} = 980$ MeV and $m_{\phi} = 1020$ MeV are used in the present work.

IV. RESULTS AND DISCUSSION

In this section, we investigate the properties of hyperonic matter in the presence of strong magnetic fields using the RMF theory. We adopt the TM1 version of the RMF theory, which is known to provide excellent descriptions of the ground states of finite nuclei including unstable nuclei [39]. For the meson-hyperon couplings, we use the values constrained by reasonable hyperon potentials that include the updated information from recent developments in hypernuclear physics. We take into account the two additional hidden-strangeness mesons, $\sigma^*$ and $\phi$, and consider two cases of the $YY$ interaction. The weak $YY$ interaction is suggested by the Nagara event, whereas the strong $YY$ interaction is deduced from the earlier measurement. By comparing the results with different $YY$ interactions, we evaluate how sensitive the results are to this factor.

It has been found in Refs. [4, 5, 6, 7, 8, 9, 10] that the properties of neutron star matter can be significantly affected only for the magnetic field strength $B^* = B/B^c_e > 10^5$ ($B^c_e = 4.414 \times 10^{13}$ G is the electron critical field). In Fig. 1, we plot the meson mean fields obtained as a function of the baryon density. The results with $B^* = 10^6$ are shown in the right panels, whereas those in the field-free case are shown in the left panels. We note that the mean fields of $\sigma$ and $\sigma^*$ mesons have negative values in this calculation. It is seen that $-\sigma$ in the right panels is smaller than the one in the left panels, but $\omega$ is a little larger. $-\sigma^*$ and $\phi$ in the right panels appear later than those in the left panels, because the onset of hyperons shifts to higher density in the presence of strong magnetic fields. By comparing the upper and lower panels, we find that $-\sigma^*$ and $\phi$ with the weak $YY$ interaction are smaller than those with the strong $YY$ interaction. This is because the hyperon couplings to $\sigma^*$ in the weak $YY$ case are smaller than those in the strong $YY$ case (see Table III). In Fig. 2
we show the effective masses of nucleons (left panels) and \( \Lambda \) hyperons (right panels) as a function of the baryon density for \( B^* = 0, 10^5, \) and \( 10^6 \). We notice that other hyperons (\( \Sigma \) and \( \Xi \)) have a similar behavior to \( \Lambda \). It is found that the influence of the magnetic field on the effective masses is not observable until \( B^* > 10^5 \). The effective masses for \( B^* = 10^6 \) are significantly larger than the field-free values. This is because smaller \(-\sigma\) and \(-\sigma^*\) are obtained in the presence of strong magnetic fields as shown in Fig. 1. Note that the effective mass of baryon species \( b \) is given by \( m^*_b = m_b + g_{\sigma b} \sigma + g_{\sigma^* b} \sigma^* \). The decrease of the effective mass with increasing density is much larger in the left panels than in the right panels, which is mainly due to \( g_{\sigma N} > g_{\sigma \Lambda} \). Because the hyperon couplings to \( \sigma^* \) in the weak \( YY \) case are smaller than those in the strong \( YY \) case (see Table 2), the effective masses of \( \Lambda \) in the upper right panel of Fig. 2 are larger than those in the lower right panel. On the other hand, the effective masses of nucleons are not affected by \( \sigma^* \) due to \( g_{\sigma^* N} = 0 \).

In Figs. 3, 4, and 5, we present the particle fraction \( Y_i = \rho_i^b / \rho_b \) as a function of the baryon density \( \rho_b \) for \( B^* = 0, 10^5, \) and \( 10^6 \), respectively. At low densities, the fractions of nucleons and leptons are significantly affected by the magnetic field. For \( B^* = 0 \) shown in Fig. 3 the proton fraction rises rapidly with increasing density and reaches \( \sim 0.2 \) around \( 2 \rho_0 \) before the appearance of hyperons. For \( B^* = 10^5 \) (see Fig. 4), the proton fraction stays at \( \sim 0.2 \) from low densities until the appearance of hyperons at \( \sim 2 \rho_0 \). The proton fraction for \( B^* = 10^6 \) (Fig. 5) is larger than 0.5 before hyperons appear at \( \sim 4 \rho_0 \). It is obvious that the threshold densities of hyperons are significantly altered by the magnetic field in Fig. 5. \( \Lambda \) is the first hyperon to appear in all cases. In Figs. 3 and 4, \( Y_\Lambda \) is found to be larger than other hyperon fractions, but it is much smaller than \( Y_{\Xi^-} \) and \( Y_{\Xi^0} \) at higher densities in Fig. 5. This is because \( \Lambda \) is a neutral particle and has a small anomalous magnetic moment. In general, charged particles depend more on the magnetic field than neutral particles due to the Landau quantization of charged particles. For neutral particles, the anomalous magnetic moment plays an important role in the presence of strong magnetic fields. We find that hyperon fractions are very sensitive to hyperon couplings used in the calculation. The appearance of hyperons leads to pronounced changes in the fractions of nucleons and leptons. The negatively charged hyperons can play the same role as electrons and muons in maintaining the electric charge neutrality. Therefore, the appearance of \( \Xi^- \) decreases \( Y_e \) and \( Y_\mu \). By comparing the upper and lower panels of Figs. 3, 4, we find that the differences between weak and strong \( YY \) cases are not very large and mainly exist at high densities.
In Fig. 6, we show the matter pressure $P$ as a function of the matter energy density $\varepsilon$ for the magnetic field strengths $B^* = 0$, $10^5$, and $10^6$. The results with the weak and strong YY interactions are plotted in the upper and lower panels, respectively. It is found that the EOS with the weak YY interaction is slightly stiffer than the one with the strong YY interaction. The presence of hyperons makes the EOS softer than in the case without hyperons, which has been discussed in our previous work [9]. The softening of the EOS becomes less pronounced with increasing magnetic field for $B^* > 10^5$. This is because the onset of hyperons shifts to higher densities, and the hyperon contribution becomes smaller with increasing $B^*$. Here we include the anomalous magnetic moments of all baryons, which play an important role in determining the EOS and composition of hyperonic matter.

V. SUMMARY

In this article, we have studied the effects of strong magnetic fields on the properties of hyperonic matter. We have employed the RMF theory with the TM1 parameter set which is known to provide excellent descriptions of nuclear matter and finite nuclei including unstable nuclei. In the RMF approach, baryons interact through the exchange of scalar and vector mesons. The baryons considered in this work are nucleons ($p$ and $n$) and hyperons ($\Lambda$, $\Sigma$, and $\Xi$). The exchanged mesons include isoscalar scalar and vector mesons ($\sigma$ and $\omega$), isovector vector meson ($\rho$), and two additional hidden-strangeness mesons ($\sigma^*$ and $\phi$). It is well known that the meson-hyperon couplings play an important role in determining the properties of hyperonic matter. We have used the couplings constrained by reasonable hyperon potentials that include the updated information from recent developments in hypernuclear physics. The weak hyperon-hyperon ($YY$) interaction suggested by the Nagara event has been adopted to investigate the properties of hyperonic matter, whereas the strong $YY$ interaction deduced from the earlier measurement has also been used for comparison. We found that the EOS with the weak $YY$ interaction is slightly stiffer than the one with the strong $YY$ interaction, and the differences mainly exist at high densities.

It is found that the effects of strong magnetic fields become significant only for magnetic field strength $B^* > 10^5$. The threshold densities of hyperons can be significantly altered by the magnetic field. The presence of hyperons makes the EOS softer than in the case without hyperons, and the softening of the EOS becomes less pronounced with increasing
We found that $\Lambda$ is the first hyperon to appear in all cases considered here. For strong magnetic fields as shown in Fig. 5, the fractions of $\Lambda$ are much smaller than those of $\Xi^-$ and $\Xi^0$ at higher densities. This is because $\Lambda$ is a neutral particle and has a small anomalous magnetic moment. It is found that charged particles depend more on the magnetic field than neutral particles due to the Landau quantization of charged particles. In this work, we have included the anomalous magnetic moments of all baryons, which play an important role in determining the EOS and composition of hyperonic matter.

Acknowledgment

This work was supported in part by the National Natural Science Foundation of China (No. 10675064).
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TABLE I: Static properties of baryons considered in this work. The mass and charge of baryon species \( b \) are denoted by \( m_b \) and \( q_b \), respectively. The baryonic magnetic moment is denoted by \( \mu_b \), and the anomalous magnetic moment is given by \( \kappa_b = (\mu_b/\mu_N - q_b m_p/m_b)\mu_N \), where \( \mu_N = e\hbar/2m_p \) is the nuclear magneton.

| \( b \) | \( m_b \) (MeV) | \( q_b \) | \( \mu_b \) (\( \mu_N \)) | \( \kappa_b \) (\( \mu_N \)) |
|--------|----------------|---------|----------------|----------------|
| \( p \)  | 938.0          | 1       | 2.79           | 1.79           |
| \( n \)  | 938.0          | 0       | -1.91          | -1.91          |
| \( \Lambda \) | 1115.7        | 0       | -0.61          | -0.61          |
| \( \Sigma^+ \) | 1193.1        | 1       | 2.46           | 1.67           |
| \( \Sigma^0 \) | 1193.1        | 0       | 1.61           | 1.61           |
| \( \Sigma^- \) | 1193.1        | -1      | -1.16          | -0.37          |
| \( \Xi^0 \)  | 1318.1         | 0       | -1.25          | -1.25          |
| \( \Xi^- \)  | 1318.1         | -1      | -0.65          | 0.06           |

TABLE II: Scalar coupling constants determined by the hyperon potentials. We take \( g_{\sigma^*\Sigma} = g_{\sigma^*\Lambda} \) in this calculation.

|          | \( g_{\sigma\Lambda} \) | \( g_{\sigma\Sigma} \) | \( g_{\sigma\Xi} \) | \( g_{\sigma^*\Lambda} \) | \( g_{\sigma^*\Xi} \) |
|---------|--------------------------|------------------------|---------------------|--------------------------|------------------------|
| Weak YY | 6.228                    | 4.472                  | 3.114               | 5.499                    | 11.655                 |
| Strong YY | 6.228                   | 4.472                  | 3.114               | 7.103                    | 12.737                 |
FIG. 1: (Color online) Meson mean fields as a function of baryon density, for $B^* = 10^6$ and the field-free ($B^* = 0$) case, with the weak and strong $YY$ interactions.

FIG. 2: (Color online) Effective masses of nucleons (N) and $\Lambda$ hyperons as a function of baryon density for $B^* = 0$, $10^5$, and $10^6$, and for weak and strong $YY$ interactions.
FIG. 3: Particle fraction \( Y_i = \rho_i^v / \rho_b \) as a function of baryon density \( \rho_b \) for \( B^* = 0 \), with weak and strong YY interactions.

FIG. 4: Same as Fig. 3 but for \( B^* = 10^5 \).
FIG. 5: Same as Fig. 3 but for $B^* = 10^6$.

FIG. 6: (Color online) Matter pressure $P$ vs the matter energy density $\varepsilon$ for $B^* = 0$, $10^5$, and $10^6$, for weak and strong $YY$ interactions.