We construct a new isentropic equation of state (EOS) at finite temperature, "CRover," on the basis of the hadron–quark crossover at high density. By using the new EOS, we study the structure of hot neutron stars at birth with typical lepton fraction ($Y_l = 0.3 - 0.4$) and typical entropy per baryon ($\hat{S} = 1 - 2$). Due to the gradual appearance of quark degrees of freedom at high density, the temperature $T$ and the baryon density $\rho$ at the center of hot neutron stars with hadron–quark crossover are found to be smaller than those without the crossover by a factor of two or more. Typical energy release due to the contraction of a hot neutron star to a cold neutron star with mass $M = 1.4 M_\odot$ is shown to be about $0.04 M_\odot$, with a spin-up rate of about 14%.

Introduction

In a core-collapse Type-II supernova explosion, a proto-neutron star (PNS) with radius $\sim 100 - 200$ km is formed. During the first few seconds after the core bounce, the PNS undergoes a rapid contraction and evolves into either a "hot" neutron star (NS) with radius $\sim 10 - 20$ km or a black hole. The hot NS at birth in quasi-hydrostatic equilibrium is composed of supernova matter with a typical lepton fraction of $Y_l = Y_e + Y_\nu \sim 0.3 - 0.4$ and a typical entropy per baryon$^1$ of $\hat{S} \sim 1 - 2$; these values are caused by neutrino trapping at a baryon density $\rho$ exceeding $10^{12}$ g cm$^{-3}$. With this as an initial condition, the hot NS contracts gradually by neutrino diffusion with a time scale of several tens of seconds and evolves to a nearly "cold" NS with $Y_\nu \simeq 0$ and $\hat{S} \simeq 0$, unless another collapse to a black hole takes place [1–3].

The hot neutron star provides us with a variety of information on the properties and dynamics of high density matter [4,5]. The purpose of this Letter is to study the hot neutron star at birth with degenerate neutrinos on the basis of an equation of state (EOS) "CRover" at finite temperature $T$ which is newly developed on the basis of the hadron–quark crossover picture at high baryon densities. Such an EOS for "cold" neutron-star matter has previously been studied by the present authors [6,7]: It was shown that a smooth crossover from hadronic matter to strongly interacting quark matter around $\rho \sim 3 \rho_0 (\rho_0 = 0.17$ fm$^{-3}$ being the normal nuclear matter density) can support

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$^1$ These authors contributed equally to this work.

$^2$ Throughout this Letter, we use a hat symbol for the thermodynamic quantities per baryon.
a cold neutron star with a maximum mass greater than the observed masses of massive neutron stars, \( M = (1.97 \pm 0.04) M_\odot \) [8] and \( M = (2.01 \pm 0.04) M_\odot \) [9]. Similar conclusions for cold neutron stars have recently been reported by several other groups [10–12].

In this Letter, we will generalize the idea of the hadron–quark crossover to a system at finite \( T \) by considering the Helmholtz free energy \( F(N, V, T) \) as a basic quantity to interpolate the hadronic matter and the quark matter. Here, \( N \) and \( V \) are the total baryon number and the total volume, respectively. By using our EOS for supernova matter with fixed \( Y_l \) and \( S \), we study not only the bulk observables such as the \( M\sim R \) relation of the hot neutron star at birth, but also the temperature, density, and sound velocity profiles inside the star. We note here that hot neutron stars with a hadron–quark mixed phase have previously been studied, e.g. [13–16]. Such a mixed phase generally leads to a soft EOS, so that it is rather difficult to sustain \( 2M_\odot \) NSs. On the contrary, our hadron–quark crossover approach does not suffer from this problem, since it leads to a stiff EOS in the crossover region.

**EOS for supernova matter at finite \( T \) with hadrons and leptons** A set of finite-temperature Hartree–Fock (HF) equations is solved for isothermal matter composed of \( n, p, e^-, e^+, \nu_e, \) and \( \bar{\nu}_e \), under the conditions of charge neutrality, chemical equilibrium, and baryon and lepton number conservations. It gives the single-particle energies, chemical potentials, and mixing ratios of nucleons and leptons, so that the hadron + lepton EOS (HL-EOS) for a given set of \( T, \rho \), and \( Y_l \) is obtained. Then, we calculate the thermodynamic quantities such as the internal energy \( E \), the entropy \( S \), and the free energy \( F(=E-TS) \)—the details are shown in [17]. We neglect hyperons \( (Y) \) in the hadronic matter in this Letter, since they do not appear as hadronic degrees of freedom once the hadron–quark crossover takes place at around \( \rho \simeq 3\rho_0 \), as shown in [6,7].

We employ \( \tilde{V}_{NN}(=\tilde{V}_{RSC}+\tilde{U}_{TNI}) \) for the effective \( NN \) interaction in the HF calculation. Here, \( \tilde{V}_{RSC} \) is a two-nucleon potential based on \( G \)-matrix calculations with the Reid soft-core potential; this depends on the total baryon density and the density difference between neutrons and protons. On the other hand, \( \tilde{U}_{TNI} \) is based on the phenomenological three-nucleon potential of Friedman–Pandharipande type [19] and is written in the form of a \( \rho \)-dependent two-body interaction. We introduce \( \tilde{U}_{TNI} \) to ensure the correct saturation of symmetric nuclear matter with nuclear incompressibility \( \kappa = 250 \text{ MeV} \) consistent with experiments. This is called the TNI2-EOS [17]. The maximum mass of the NS at \( T=0 \) with the TNI2-EOS alone is 1.63 \( M_\odot \). However, the hadron–quark crossover inside the NSs can bring this well above \( 2M_\odot \), as shown in [6,7].

**EOS for supernova matter at finite \( T \) with quarks and leptons** To obtain the quark + lepton EOS (QL-EOS), we follow [6,7] and apply the \((2+1)\)-flavor Nambu–Jona-Lasinio (NJL) model,

\[
\mathcal{L}_{NJL} = \overline{q}(i\gamma \partial - m)q + \frac{G_S}{2} \sum_{a=0}^{8} \left( (\overline{q} \gamma^a q)^2 + (\overline{q} i \gamma_5 \lambda^a q)^2 \right) \\
- G_D \left[ \det \overline{q} (1 + \gamma_5) q + \text{h.c.} \right] - \frac{G_V}{2} (\overline{q} \gamma^\mu q)^2. \tag{1}
\]

Here, the quarks \( q_i \) \( (i = u, d, s) \) with current quark masses \( m_i \) have three colors and three flavors. The term proportional to \( G_S \) in Eq. \( (1) \) is a \( U(3)_L \times U(3)_R \) symmetric four-fermi interaction where \( \lambda^a \) are the Gell-Mann matrices with \( \lambda^0 = \sqrt{2/3}I_1 \). The term proportional to \( G_D \) is the Kobayashi–Maskawa–’t Hooft (KMT) six-fermi interaction which breaks \( U(1)_A \) symmetry explicitly. The last
term with the positive coupling $g_V$ represents a universal repulsion among different flavors. A standard parameter set obtained from hadron phenomenology in the vacuum [18] reads $\Lambda$ (UV cutoff) = 631.4 MeV, $G_s \Lambda^2 = 3.67$, $G_D \Lambda^3 = 9.29$, $m_{u,d} = 5.5$ MeV, and $m_s = 135.7$ MeV. As for the universal vector coupling, we consider the typical value, $g_V / G_S = 0.5$. The maximum mass of a cold NS with hadron–quark crossover with these parameters reads 2.6 $M_\odot$ (see Fig. 19 and Table 7 of [7]).

The Gibbs free energy $\Omega$ is the sum of the quark contribution and the lepton contribution, $\Omega(\mu, V, T; \mu_i) = \Omega_{\text{quark}}(\mu, V, T; \mu_i) + \Omega_{\text{lepton}}(\mu, V, T; \mu_i)$, with $\mu$ and $\mu_i$ being the baryon chemical potential and the lepton chemical potential respectively.\(^2\) The former in the mean field approximation is

\[
\frac{\Omega_{\text{quark}}}{V} = -T \sum_i \sum_\ell \int \frac{d^3 p}{(2\pi)^3} \text{Tr} \ln \left( \frac{S_i^{-1}(\omega_\ell, p)}{T} \right) + G_S \sum_i \sigma_i^2 + 4 G_D \sigma_u \sigma_d \sigma_s - \frac{1}{2} g_V \left( \sum_i n_i \right)^2,
\]

where $n_i = \{q_i, \bar{q}_i\}$ is the quark number density in each flavor ($i = u, d, s$). $S_i$ is the quark propagator, which can be written as

\[
S_i^{-1}(\omega_\ell, p) = \not{p} - M_i - \gamma^0 \mu_i^{\text{eff}}, \quad \mu_i^{\text{eff}} \equiv \mu_i - g_V \sum_j n_j,
\]

where $\omega_\ell = (2 \ell + 1) \pi T$ is the Matsubara frequency and $\mu_i^{\text{eff}}$ is the effective chemical potential. The lepton part $\Omega_{\text{lepton}}(\mu, V, T; \mu_i)$ corresponds to a relativistic and degenerate system of electrons, neutrinos, and their anti-particles. Charge neutrality and chemical equilibrium among quarks and leptons are imposed. The Helmholtz free energy $F(N, V, T; Y_l)$ is obtained from the Gibbs free energy $\Omega(\mu, V, T; \mu_i)$ by the Legendre transformation.

**Hadron–quark crossover at finite T** Let us now generalize the idea of the hadron–quark crossover introduced in our previous works [6,7] to the system at finite $T$. We consider the following interpolation of the Helmholtz free energy per baryon $\hat{F}$ between the hadron + lepton free energy per baryon ($\hat{F}_{\text{HL}}$) and the quark + lepton free energy per baryon ($\hat{F}_{\text{QL}}$),

\[
\hat{F}(\rho, T; Y_l) = \hat{F}_{\text{HL}}(\rho, T; Y_l) w_-(\rho, T) + \hat{F}_{\text{QL}}(\rho, T; Y_l) w_+(\rho, T).
\]

Here, $w_-$ and $w_+ = 1 - w_-$ are the weight functions. The typical temperature we consider is 30 MeV or less. This is considerably smaller than the thermal dissociation temperature ($\sim$200 MeV) of hadrons into quarks and gluons [20], so we assume that $w_-$ is $T$-independent and takes the same form given in [6,7]:

\[
w_\pm(\rho, T) \rightarrow w_\pm(\rho) = \frac{1}{2} \left( 1 \pm \tanh \left( \frac{\rho - \bar{\rho}}{\Gamma} \right) \right),
\]

where $\bar{\rho}$ and $\Gamma$ are the phenomenological interpolation parameters which characterize the crossover density and the width of the crossover window, respectively. We take $(\bar{\rho}, \Gamma) = (3 \rho_0, \rho_0)$ as typical values to account for 2 $M_\odot$ neutron stars at $T = 0$ [6,7]. Note that hyperons do not appear in the hadronic phase with such a relatively low crossover density.

\(^2\) The thermal contribution from massless gluons above the crossover density is less than 4% of the thermal part of quarks for $T < 30$ MeV, so we do not consider gluons in this Letter.
From Eqs. (4) and (5), other bulk quantities, the entropy per baryon (\(\hat{S}\)) and the energy per baryon (\(\hat{E}\)), are derived by using the thermodynamic relations \(\hat{S} = -\partial \hat{F} / \partial T\) and \(\hat{E} = \hat{F} + T \hat{S}\) as

\[
\hat{S}(\rho, T; Y_l) = \hat{S}_{\text{HL}}(\rho, T; Y_l)w_-(\rho) + \hat{S}_{\text{QL}}(\rho, T; Y_l)w_+(\rho),
\]

\[
\hat{E}(\rho, T; Y_l) = \hat{E}_{\text{HL}}(\rho, T; Y_l)w_-(\rho) + \hat{E}_{\text{QL}}(\rho, T; Y_l)w_+(\rho).
\]

The thermodynamic quantities for isothermal matter with \(\rho, T\), and \(Y_l\) can be converted into those for isentropic matter by means of the \(T-\rho\) relationship constrained by the constant entropy per baryon, \(\hat{S} = \text{const.}\) Then the isentropic pressure \(P\) and the isentropic baryon chemical potential \(\mu_b\) are obtained as

\[
P\left(\rho, T(\rho); Y_l, \hat{S}\right) = -\frac{\partial E}{\partial V}_{S,N} = \rho^2 \frac{\partial \hat{E}}{\partial \rho}_{\hat{S}},
\]

\[
\mu_b\left(\rho, T(\rho); Y_l, \hat{S}\right) = \frac{\partial E}{\partial N}_{S,V} = \frac{\partial e}{\partial \rho}_{\hat{S}},
\]

where \(e \equiv \rho \hat{E}\) is the energy density. From the definition of our interpolation method Eq. (4), \(P\) and \(\mu_b\) approach those of pure hadronic (quark) matter in the low (high) density limit.

**EOS for supernova matter below \(\rho_0\)**  Below the normal nuclear matter density \(\rho_0\), we use the thermal EOS which consists of an ensemble of nuclei and interacting nucleons in nuclear statistical equilibrium by Hempel and Schaffner-Bielich [21], which we call the HS EOS in the following. As the density decreases, the temperature also decreases monotonically; for example, \(T \sim 10\) (2) MeV at \(\rho = 0.1\) (0.001) fm\(^{-3}\) for \((Y_l, \hat{S}) = (0.3, 1)\). Other EOSs in this region do not show a quantitative difference from the HS EOS, as discussed in [22]. Once the baryon density becomes smaller than the neutron drip density \(10^{-3}\rho_0\) the temperature becomes smaller than 1 MeV, so that we switch to the standard BPS EOS [23]. In Fig. 1, the internal structures of hot and cold NSs with our isentropic EOS are illustrated.

**EOS for supernova matter with hadron–quark crossover**  To carry out the conversion from the isothermal matter with fixed \(T\) and \(Y_l\) to the isentropic matter with fixed \(\hat{S}\) and \(Y_l\), we consider \(\hat{S}(\rho, T, Y_l)\) as a function of \(\rho/\rho_0\). In Fig. 2(a), we show \(\hat{S}\) for \(T = 10, 15, 20\) MeV and \(Y_l = 0.3\) with crossover by the solid lines, while \(\hat{S}\) for the pure hadronic matter and the pure quark matter are shown by the dashed lines. One finds that the entropy per baryon turns out to increase at high densities, once the quark degrees of freedom start to appear. Figure 2(b) shows the temperature as a function of \(\rho/\rho_0\) under isentropic conditions with crossover (solid lines) for \(\hat{S} = 1, 2\) and \(Y_l = 0.3\) and without crossover (the dashed line) for \(\hat{S} = 2\) and \(Y_l = 0.3\). The rapid increase of \(\hat{S}\) given \(T\) in the crossover region as shown in Fig. 2(a) can be understood as follows. First of all, the temperature is smaller than the Fermi energy at high densities, enough so that the low \(T\) expansion (the Sommerfeld expansion) is applicable. Then let us consider, for illustrative purposes, non-relativistic neutron matter in the hadronic phase and massless u–d quark matter with \(\beta\) equilibrium in the quark phase. The ratio of the entropy per baryon of the hadronic matter \((\hat{S}_H)\) and that of the quark matter \((\hat{S}_Q)\) for the non-interacting case can be evaluated in the leading order of the Sommerfeld expansion as
Fig. 1. Schematic illustrations of internal structure and relevant EOS in hot and cold NSs with $M = 2 M_\odot$ obtained by a new EOS “CRover” with hadron–quark crossover at finite temperature. We take $\left(Y_l, \hat{S}\right) = (0.3, 1)$ for the hot NS in this figure. BPS, HS, TNI2, and NJL are the EOS adopted (details are in the text).

$$\frac{\hat{S}_Q}{\hat{S}_H} \sim 1.14 \frac{g_Q^{1/3}}{g_H^{1/3}} \left(\frac{\rho}{\rho_0}\right)^{1/3}.$$ Here, the fractional powers originate from the difference between relativistic and non-relativistic kinematics. Note that the right-hand side is $T$ independent. By taking $g_Q = 2 \text{spin} \times 3 \text{color}$ and $g_H = 2 \text{spin}$, as well as the typical crossover density $\rho = 3\rho_0$, the ratio becomes 1.88. This essentially explains the jump between the two dashed lines in Fig. 2(a) for given $T$. Because of the fact that $\hat{S}$ given $T$ with the crossover [the solid line in Fig. 2(a)] becomes larger than $\hat{S}_H$ without the crossover, there is no need to increase $T$ to keep $\hat{S}$ constant at high densities. This can indeed be seen from Fig. 2(b), in which $T$ is rather insensitive to $\rho$ for fixed $\hat{S}$.

In Fig. 3(a), (b), (c), and (d), we show the isentropic pressure $P(\rho, T(\rho), Y_l)$, the energy per baryon $E(\rho, T(\rho), Y_l)$, the baryon chemical potential $\mu_B$ as a function of $\rho$, and the $P-e$ relation, respectively, for three characteristic sets, $\left(Y_l, \hat{S}\right) = (0.3, 1), (0.3, 2)$, and $(0.4, 1)$. For comparison, the EOS of cold neutron star matter ($T = 0$ without neutrino degeneracy) is also shown by the black solid lines. We call this new EOS with hadron–quark crossover “CRover,” and will post the numerical tables for different combinations of $\hat{S}$ and $Y_l$ in the EOS database (EOSDB) [24].

http://asphwww.ph.noda.tus.ac.jp/eos-gate/.
Structure of hot neutron stars

We now solve the Tolman–Oppenheimer–Volkov (TOV) equation to obtain the structure of a hot neutron star at birth:

\[
\frac{dP}{dr} = -\frac{G}{r^2} \left( M(r) + 4\pi Pr^3 \right) \left( \varepsilon + P \right) \left( 1 - 2GM(r)/r \right)^{-1},
\]

\[
M(r) = \int_0^r 4\pi r'^2 \varepsilon(r') dr'.
\]

(10)

Here we have assumed spherically symmetric stars with \( r \) being the radial distance from the center. In the following, we will consider hot neutron stars with typical values, \( (Y_1, \hat{S}) = (0.3, 1) \).

In Fig. 4(a), we show the relation between the gravitational mass \( M \) and the baryon number \( N_B \) for hot and cold neutron stars with and without crossover. First of all, the crossover leads to heavier neutron stars for given \( N_B \) due to the stiffening of the EOS. Secondly, hot neutron stars have larger mass than the cold ones for given \( N_B \). Moreover, the maximum \( N_B \) for a hot NS is smaller than that of a cold NS.\(^4\) Typical energy release for \( M_{\text{cold}} = 1.4 M_\odot \) due to the contraction reads \( \Delta E = M_{\text{hot}} - M_{\text{cold}} \sim 0.04 M_\odot \). In Fig. 4(b), we show \( \Delta E \) in units of \( M_\odot \) as a function of the mass of a cold neutron star, \( M_{\text{cold}} \).

In Fig. 5(a), the mass–radius (\( M–R \)) relations of hot and cold neutron stars with hadron–quark crossover are shown. The stiffening of the EOS at finite \( T \) at low density [see Fig. 3(d)] makes the hot neutron star bigger in size, especially for neutron stars with small \( M \). On the other hand, a slight

\(^4\) With the present hadron–quark crossover, the hot EOS is slightly softer than the cold EOS in the crossover region as shown in Fig. 3(d), so that the delayed collapse of hot NS to black hole does not take place. On the other hand, with the first-order hadron–quark phase transition, the softening of the EOS due to the mixed phase is tamed by the finite temperature effect, so that hot NSs may eventually collapse into black holes [13,15]. The possibility of delayed collapse was reported earlier by one of the present authors in the context of pion condensation [4].
Fig. 3. (a) The isentropic pressure $P$ as a function of baryon density $\rho$ for $(Y_t, \hat{S}) = (0.3, 1), (0.3, 2)$, and $(0.4, 1)$. The black line corresponds to the EOS for cold neutron star matter. The crossover window is shown by the shaded area on the horizontal axis. (b) The energy per baryon $\hat{E}$ with the same set of $Y_t$ and $\hat{S}$ as (a). (c) Baryon chemical potential $\mu_B$ as a function of $\rho$. (d) $P$ as a function of $\varepsilon$.

softening of the EOS at finite $T$ in the crossover region [see Fig. 3(d)] leads to the maximum mass of the hot neutron star being slightly smaller than that of the cold neutron star. The main reason behind such a slight softening of EOS is the decrease of $T$ around the crossover region as shown in Fig. 2(b). Further studies would be necessary, however, to check how this softening depends on the weight function $w_{\pm}$. The sound velocity of the system, which is one of the measures of the stiffness of the EOS, is obtained from

$$v_s^2(\rho; Y_t, \hat{S}) = \frac{\partial P}{\partial \varepsilon}_{Y_t, \hat{S}} = \frac{dP(\rho, T(\rho); Y_t, \hat{S})/d\rho}{d\varepsilon(\rho, T(\rho); Y_t, \hat{S})/d\rho}_{Y_t, \hat{S}}.$$  \hspace{1cm} \text{(11)}$$

As we have demonstrated at $T = 0$ in [7], $v_s^2$ becomes large in the crossover region. We find the same tendency at finite $T$ too. In Fig. 5(b), we plot the local sound velocity squared $v_s^2(r)$ as a function of $r$ for $M = 1.4M_\odot$ obtained by Eq. (11) with $\rho(r)$ obtained by the TOV equation. We
Fig. 4. (a) The neutron star mass $M$ as a function of the total baryon number $N_B$. Red (blue) curves correspond to the hot (cold) neutron stars with crossover. The dotted lines correspond to the case without crossover. (b) The energy release $\Delta E$ as a function of the cold neutron star mass $M_{\text{cold}}$. $(Y_f, \hat{S}) = (0.3, 1)$ is adopted.

note that the sound velocity decreases toward the neutron star surface simply because the baryon density decreases.

In Fig. 6(a) and (b), the central temperature $T_{\text{cent}}$ and the central density $\rho_{\text{cent}}$ of hot neutron stars with and without the crossover are plotted as a function of $M$. The temperature decreases due to the appearance of the quark degrees of freedom [see Fig. 2(b)], so that the central temperature becomes significantly smaller with crossover, as shown in Fig. 6(a). The central density of the star becomes significantly smaller with crossover, as shown Fig. 6(b), due to the fact that the EOS is stiffer with crossover.

To compare the internal structure of the hot neutron star with and without the crossover in terms of the temperature and density, we plot $T(r)$ and $\rho(r)$ in Fig. 7 (a) and (b), respectively. Here we consider a hot neutron star with a canonical mass $M = 1.4 M_\odot$. In the presence of the hadron–quark crossover, the EOS is stiffer and hence the central density is smaller, so that the temperature and density profiles of the star are more uniform as compared to those without the crossover.
Fig. 5. Mass $(M)$-radius $(R)$ relationship for $(Y_l, \hat{S}) = (0.3, 1)$. Red: hot neutron stars with crossover. Blue: cold neutron stars with crossover. (b) The sound velocity squared $v_s^2$ as a function of the distance $r$ from the center of a $1.4 M_\odot$ neutron star. The colors of each line are the same as in (a).

Fig. 6. (a) The central temperature $T_{\text{cent}}$ as a function of $M$. (b) The central density $\rho_{\text{cent}}$ as a function of the neutron star mass $M$. The solid and dashed lines correspond to the EOS with (without) crossover, and $(Y_l, \hat{S}) = (0.3, 1)$ is taken.

Summary and concluding remarks In this Letter, we have discussed the properties of hot neutron stars at birth on the basis of a new EOS “CRover” for supernova matter with hadron–quark crossover. Such a crossover leads to an EOS stiff enough to sustain $2 M_\odot$ neutron stars. A noticeable point is that the crossover plays important roles not only in generating a stiff EOS but also in lowering the internal temperature of hot neutron stars. Such suppression of temperature originates from a combined effect of the isentropic nature of the supernova matter and larger entropy for a given temperature due to the quark degrees of freedom. For a given baryon number, hot neutron stars have larger radius and larger gravitational mass caused by the high lepton fraction and the thermal effect. This suggests that, during the contraction from hot to cold stars, gravitational energy is released and simultaneously spin-up.
Fig. 7. (a) The temperature profiles of a hot neutron star with $M = 1.4 M_\odot$ and $(Y_l, \hat{S}) = (0.3, 1)$. Solid and dashed lines correspond to the EOS with crossover and without crossover respectively. (b) The density profiles of the same neutron star as case (a).

mass accretion after the birth. In the present study, we have used a temperature-independent interpolation function to construct the EOS with crossover. Although this is justified for supernova matter with temperature less than 30 MeV, further investigations would be necessary for the EOS to be applicable to a wider range of temperatures in supernovae and neutron star mergers.

Acknowledgements

We thank Gordon Baym for helpful discussions in the early stage of this work. KM thanks Mark Alford for kind hospitality and discussions at Washington University in St. Louis where part of this work was carried out under the support of the ALPS Program, University of Tokyo. KM is supported by a JSPS Research Fellowship for Young Scientists. TH and TT were partially supported by JSPS Grant-in-Aid for Scientific Research, No.25287066. This work was partially supported by the RIKEN iTHES Project.

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