Phase transition in warm nuclear matter with alternative derivative coupling models

M. Malheiro,[1,*] A. Delfino,[2] and C. T. Coelho[2]
[1]Department of Physics, University of Maryland, College Park, Maryland 20742-4111
[2]Instituto de Física, Universidade Federal Fluminense, 24210-340, Niterói, R. J., Brazil
(Received 29 December 1997)

An analysis is performed of the liquid-gas phase transition of nuclear matter obtained from different versions of scalar derivate coupling suggested by Zimanyi and Moszkowski (ZM) and the results are compared with those obtained from the Walecka model. We present the phase diagram for the models and one of them, the ZM3 model, has the lowest critical temperature $T_c = 13.6$ MeV with the lowest critical density $\rho_c = 0.037\ \mathrm{fm}^{-3}$ and pressure $p_c = 0.157\ \mathrm{MeV\ fm}^{-3}$. These results are in accordance with recent observations from energetic heavy-ion collisions, which suggest a small liquid-gas phase region. [S0556-2813(98)02807-6]

PACS number(s): 21.65.+f, 05.70.Fh, 21.30.Fe, 25.75.−q

I. INTRODUCTION

Nowadays, the study of the liquid-gas phase transition, which may occur in the warm and dilute matter produced in energetic heavy-ion collisions, is one of the most interesting problems in nuclear physics [1]. This idea, that nuclear systems may show a critical behavior, was initiated more than ten years ago with the observation by the Purdue-Fermilab group of asymptotic fragment charge distributions exhibiting a power law [2]. This interest has increased recently with the attempt by the EOS Collaboration to extract critical exponents of fragmenting nuclear systems produced in the collision of 1 GeV/nucleon Au nuclei with a carbon target [3] and with the extraction by the ALADIN/LAND Collaboration of a caloric curve resulting from the fragmentation of the quasiprojectile formed in the collision Au + Au at 600 MeV/nucleon exhibiting a behavior expected for a first-order liquid-gas phase transition [4].

At the time when the search for signals of the liquid-gas nuclear phase transition is taking place, it is important to have ready the theoretical phase-transition predictions for a broad class of different hadronic models. In fact, the liquid-gas phase transition has recently been studied by taking into account different effects such as discontinuity in the freeze-out density [5], excluded volume that suppresses the particle number densities [6], and the inclusion of a dilaton field associated with broken scale invariance which allows one to lower the compressibility [7].

The main ingredient in these analyses is the nuclear matter equation of state (EOS) at finite temperature. The success of relativistic mean-field theories describing cold nuclear matter and bulk nuclear properties throughout the periodic table suggests the use of a relativistic mean-field EOS. Moreover, the mean-field approximation is known to be thermodynamically consistent [8,9].

Recently variants of the Zimanyi-Moszkowski (ZM) model [10] were implemented and applied to dense and cold nuclear matter [11]. The aim of this paper is to extend our study to include temperature effects and to perform an analysis of the liquid-gas phase transition of the warm nuclear matter obtained from these ZM models and compare them to the linear Walecka model.

The behavior of the nuclear matter with density and temperature is also vital to describe very different astrophysics phenomena such as supernova explosions and neutron star properties. Then, a complete thermodynamic study of this modified ZM versions is really needed and a recent application of them has already been performed in a study of the density and temperature dependences of hadron masses [13].

The usual ZM model, also referred to in the literature as the derivative scalar coupling (DSC) model, consists of a derivative coupling between nucleons and scalar mesons σ. The model has been extended to include a nonlinear interaction between the nucleon and the vector meson ω. Two types of this interaction were employed and the resulting models were denoted ZM2 and ZM3. These models were designed to cure the defects of the Walecka model [12], namely, the low effective nucleon mass and the large incompressibility of nuclear matter. Each one of them is very simple since they have only two free parameters, the scalar (vector) coupling constants C_σ^2 (C_ω^2), adjusted to reproduce the binding energy (E_b) of the nuclear matter at $\rho = \rho_0$. The degrees of freedom are baryon fields (ψ), scalar meson fields (σ), and vector meson fields (ω).

In all ZM models, there are nonlinear interaction terms which in an approximate way incorporate the effect of many-body forces. After an appropriate rescaling of the Lagrangians, these models can be understood as generalizations of the Walecka model where the scalar and vector meson couplings become effectively density dependent [14]. This fact underlies a recent approach, known as the relativistic density-dependent Hartree-Fock approach [15–17], which describes finite nuclei and nuclear matter saturation properties using coupling constants that are fitted, at each density value, to the relativistic Brueckner-Hartree-Fock self-energy terms. The good agreement obtained for the ground-state properties of spherical nuclei lends support to this sort of description involving density-dependent coupling constants. Recently, a finite nuclei calculation has been performed in the ZM models, and the energy levels and ground-state properties of ^{16}O, ^{40}Ca, ^{48}Ca, ^{90}Zr, and ^{208}Pb are in good agreement with the experimental results [18]. One of

*Permanent address: Instituto de Física, Universidade Federal Fluminense, 24210-340, Niterói, R. J., Brazil.

TABLE I. Coupling constants C_σ^2 and C_ω^2, binding energy E_b (MeV) at equilibrium density ρ_o (fm^{-3}), m^*, and the incompressibility K for the indicated models.

Models	C_σ^2	C_ω^2	E_b	ρ_o	m^*	K
Walecka	357.4	273.8	−15.75	0.148	0.54	550.82
ZM	169.2	59.1	−15.90	0.160	0.85	224.71
ZM2	219.3	100.5	−15.77	0.152	0.82	198.32
ZM3	443.3	305.5	−15.76	0.149	0.72	155.74

the main conclusions of this analysis is that a modified version of the model, referred to in this paper as the ZM3 model, improves upon the original ZM model regarding the energy splitting of levels due to the spin-orbit interaction. The behavior of the quark and gluon condensates in a medium calculated in these models is also controlled by the EOS [19] and it was shown that chiral symmetry restoration requires the meson-nucleon coupling to be density dependent [20].

The original ZM model has already been applied to investigate some thermodynamic properties of nuclear matter [21]. They have calculated the isotherms, the pressure p as a function of the density ρ for different temperatures, and obtained the critical temperature which is a little bit smaller than that obtained in the Walecka model but at almost the same critical density. Then we can say that these two models have almost the same liquid-gas phase transition. They have also concluded that the nuclear matter incompressibility decreases when the temperature increases and as in the zero-temperature case the ZM model gives a softer EOS of nuclear matter at finite temperature than the Walecka model.

In this paper we present a thermodynamic analysis for these two new variants of the ZM model and obtain the effective nucleon mass, energy per nucleon, pressure, and entropy density as a function of the baryonic density at different temperatures. All these ZM models are softer and among them ZM3 is the softest. We show the isotherms, construct the phase diagram with the phase coexistence boundary, and present the critical and flash temperatures for the models. We found that the main difference between the thermodynamic of these new models, which incorporates a nonlinear interaction between the nucleon and vector meson, and the original ZM and the linear Walecka models, which do not have that interaction, is that they present a much smaller phase coexistence region. Then we conclude that this new interaction, which is stronger in the ZM3 model and is responsible at zero temperature for a much better nuclear matter phenomenology [11], is also very important at finite temperature. It produces for the ZM3 model the smallest phase coexistence region, with the lowest critical temperature, density, and pressure, which is in accordance with a small liquid-gas phase region supported by the recent experimental results.

The outline of the paper is as follows: in the next section we present the EOS at finite temperature. Section III includes our results and discussion of the thermodynamic properties of nuclear matter. Finally, we summarize.

II. NUCLEAR MATTER EOS AT FINITE TEMPERATURE

Since the models we are dealing with were discussed in detail in Refs. [10,11], here we will only present the La-

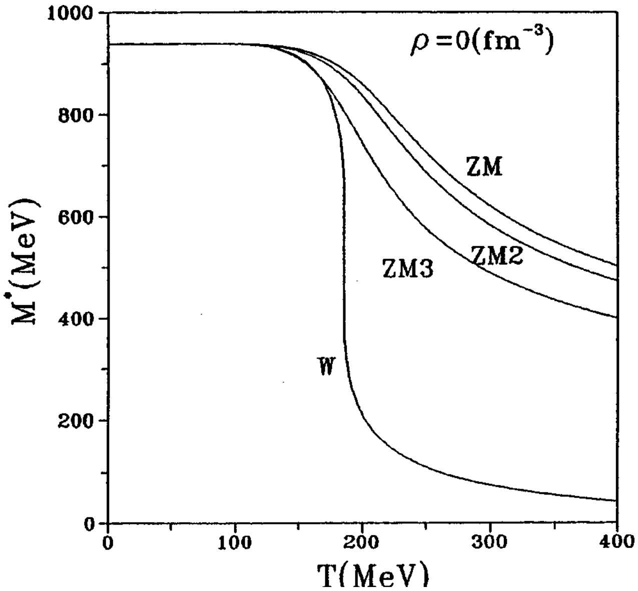

FIG. 1. Baryon effective mass in nuclear matter as a function of the temperature at $\rho=0$.

grangian obtained after rescaling the nucleon field as $\psi \to m^{*1/2}\psi$ for all ZM models and making the rescaling $\omega_\mu \to m^*\omega_\mu$ for ZM2 and ZM3 models:

$$\mathcal{L}_R=\bar{\psi}i\gamma_\mu\partial^\mu\psi+m^{*\alpha}\left(-g_\omega\bar{\psi}\gamma_\mu\psi\omega^\mu-\frac{1}{4}F^{\mu\nu}F_{\mu\nu}+\frac{1}{2}m_\omega^2\omega_\mu\omega^\mu\right)-\bar{\psi}(M-m^{*\beta}g_\sigma\sigma)\psi+\frac{1}{2}(\partial_\mu\sigma\partial^\mu\sigma-m_\sigma^2\sigma^2), \quad (1)$$

where α and β have the following values for the different models: Walecka, $\alpha=0$, $\beta=0$; ZM, $\alpha=0$, $\beta=1$; ZM2, $\alpha=1$, $\beta=1$; ZM3, $\alpha=2$, $\beta=1$; and $m^*=(1+g_\sigma\sigma/M)^{-1}$ in all three cases. M is the bare nucleon mass and $F_{\mu\nu}=\partial_\mu\omega_\nu-\partial_\nu\omega_\mu$.

When the meson fields are replaced by the constant classical fields σ_0 and ω_0 we arrive at the mean-field approximation, with the equation of motion for the nucleon:

$$[i\gamma_\mu\partial^\mu-(M-m^{*\beta}g_\sigma\sigma)-m^{*\alpha}g_\omega\gamma_\mu\omega^\mu]\psi=0, \quad (2)$$

where the effective nucleon mass M^* is given by $M^*=M-m^{*\beta}g_\sigma\sigma$. In the case of ZM models where $\beta=1$ we can identify $m^*=M^*/M=(1+g_\sigma\sigma/M)^{-1}$.

The expression for the energy density and pressure at a given temperature T can be found as usual by the average of the energy-momentum tensor,

$$\mathcal{E}=\frac{C_\omega^2}{2M^2}m^{*\alpha}\rho^2+\frac{M^4}{2C_\sigma^2}\left(\frac{1-m^*}{m^{*\beta}}\right)^2+\frac{\gamma}{(2\pi)^3}\int d^3kE^*(k)(n_k+\bar{n}_k), \quad (3)$$

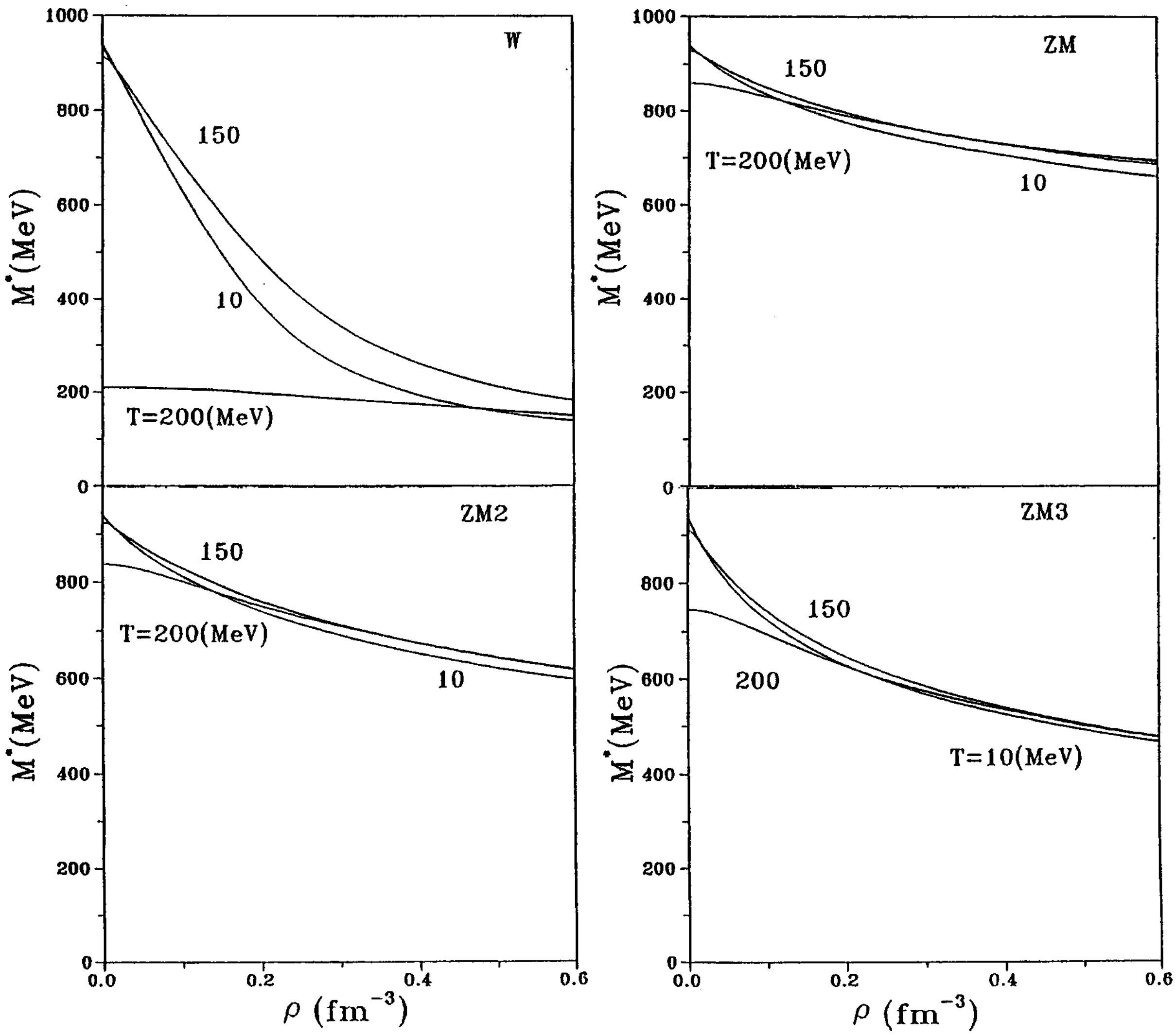

FIG. 2. Baryon effective mass M^* as a function of the baryon density at different temperatures for the Walecka model (W) and Zimanyi-Moszkowski models (ZM, ZM2, ZM3).

$$p=\frac{C_\omega^2}{2M^2}m^{*\alpha}\rho^2-\frac{M^4}{2C_\sigma^2}\left(\frac{1-m^*}{m^{*\beta}}\right)^2+\frac{1}{3}\frac{\gamma}{(2\pi)^3}\int d^3k\,\frac{k^2}{E^*(k)}(n_k+\overline{n}_k). \tag{4}$$

Thus we obtain the entropy density

$$s=\frac{1}{T}\left[\frac{C_\omega^2}{M^2}m^{*\alpha}\rho^2+\frac{\gamma}{(2\pi)^3}\int d^3kE^*(k)(n_k+\overline{n}_k)\right]+\frac{1}{3T}\frac{\gamma}{(2\pi)^3}\int d^3k\,\frac{k^2}{E^*(k)}(n_k+\overline{n}_k)-\frac{\mu\rho}{T}, \tag{5}$$

where γ is the degeneracy factor ($\gamma=4$ for nuclear matter and $\gamma=2$ for pure neutron matter), n_k and $\overline{n}_k$ stand for the Fermi-Dirac distribution for baryons and antibaryons respectively, with arguments $(E^*-\nu)/T$, and $E^*(k)$ is given by $E^*(k)=(k^2+M^{*2})^{1/2}$. An effective chemical potential which preserves the number of baryons and antibaryons in the ensemble is defined by $\nu=\mu-V$, in which μ is the thermodynamical chemical potential. We have introduced $C_\sigma^2=g_\sigma^2M^2/m_\sigma^2$ and $C_\omega^2=g_\omega^2M^2/m_\omega^2$.

The effective mass is obtained explicitly through the minimization of $\mathcal{E}$ with respect to m^* and must satisfy the self-consistent equation

$$1-m^*-\frac{\gamma C_\sigma^2}{2\pi^2}m^{*3\beta+1}\int\frac{x^2dx}{\sqrt{x^2+m^{*2}}}(n_x+\overline{n}_x)-\frac{\alpha}{2}\frac{C_\sigma^2C_\omega^2}{M^6}m^{*\alpha+2\beta}\rho^2=0, \tag{6}$$

where we have used the dimensionless variable $x=k/M$.

The energy density can be fitted to the nuclear-matter ground-state energy and saturation density ρ_0 at zero temperature to obtain the different coupling constants for the models. Only for historical reasons, since the early calculations were done in these models for $T=0$, have we maintained the values of these coupling constants in order to compare the previous results with the new ones obtained in this work at finite temperature. We would like to clarify that we could have obtained all the new coupling constants to fit the

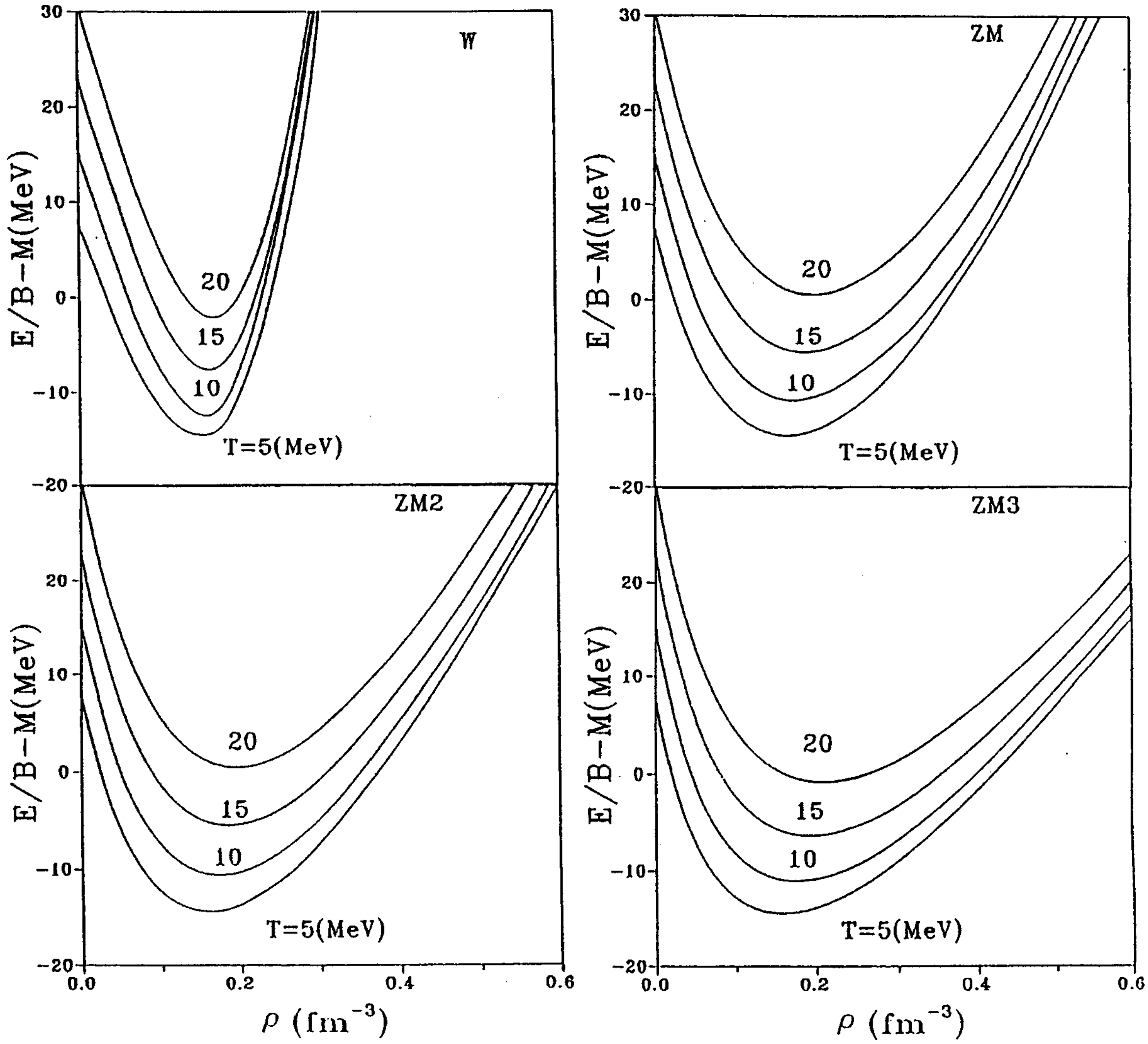

FIG. 3. Proper energy/baryon as a function of baryon density at different temperatures for the Walecka model (W) and Zimanyi-Moszkowski models (ZM, ZM2, ZM3).

same saturation point $E_b = -15.75$ MeV and $\rho_0 = 0.15\ \mathrm{fm}^{-3}$. However, these modifications will only lead to a very small change to our results and it would not affect our conclusions.

The different coupling constants for the models are presented in Table I together with the nuclear matter incompressibility which at $T=0$ is given by

$$K = 9\rho_0^2 \frac{\partial^2}{\partial \rho^2}\left(\frac{\mathcal{E}}{\rho}\right)\bigg|_{\rho=\rho_0} = 9\rho_0 \frac{\partial^2 \mathcal{E}}{\partial \rho^2}\bigg|_{\rho=\rho_0}. \tag{7}$$

The thermodynamic functions are obtained by first solving the self-consistency condition in Eq. (6) to determine M^* for each T and ν fixed. We substitute the value of M^* in the Fermi-Dirac distribution functions and then calculate the integrals in Eqs. (3), (4), and (5).

III. RESULTS AND DISCUSSION

In Fig. 1 we show M^* as a function of T at zero density. In this regime, the vector field proportional to ρ vanishes, and so the three ZM models differ only in having different values of their scalar coupling constants C_σ^2. The ZM and the Walecka models coincide in the lower-temperature region $T \leqslant 120$ MeV and the ZM3 model stays together with the Walecka model up to $T \sim 160$ MeV. However, at a higher temperature the models separate quite clearly, with the effective nucleon mass in the ZM models dropping more slowly than that in the Walecka model. This means that the sigma field (the source for the scalar density) increases more slowly with temperature in the ZM models because of the inclusion of nonlinear interactions which are absent in the Walecka model. As a result, the attraction is stronger in the Walecka model, favoring the formation of nucleon-antinucleon pairs at high temperature. Moreover, none of the proposed ZM models are able to give a first-order phase transition at $\rho=0$, $T \neq 0$. This is in contrast to the Walecka model, which has such a phase transition at $T \sim 185$ MeV [22].

In Fig. 2 we show the behavior of the effective nucleon mass with density at different temperatures for all the models. For low temperatures the results are not so different from those obtained at zero temperature, showing that in this regime the density dependence is more important than the tem-

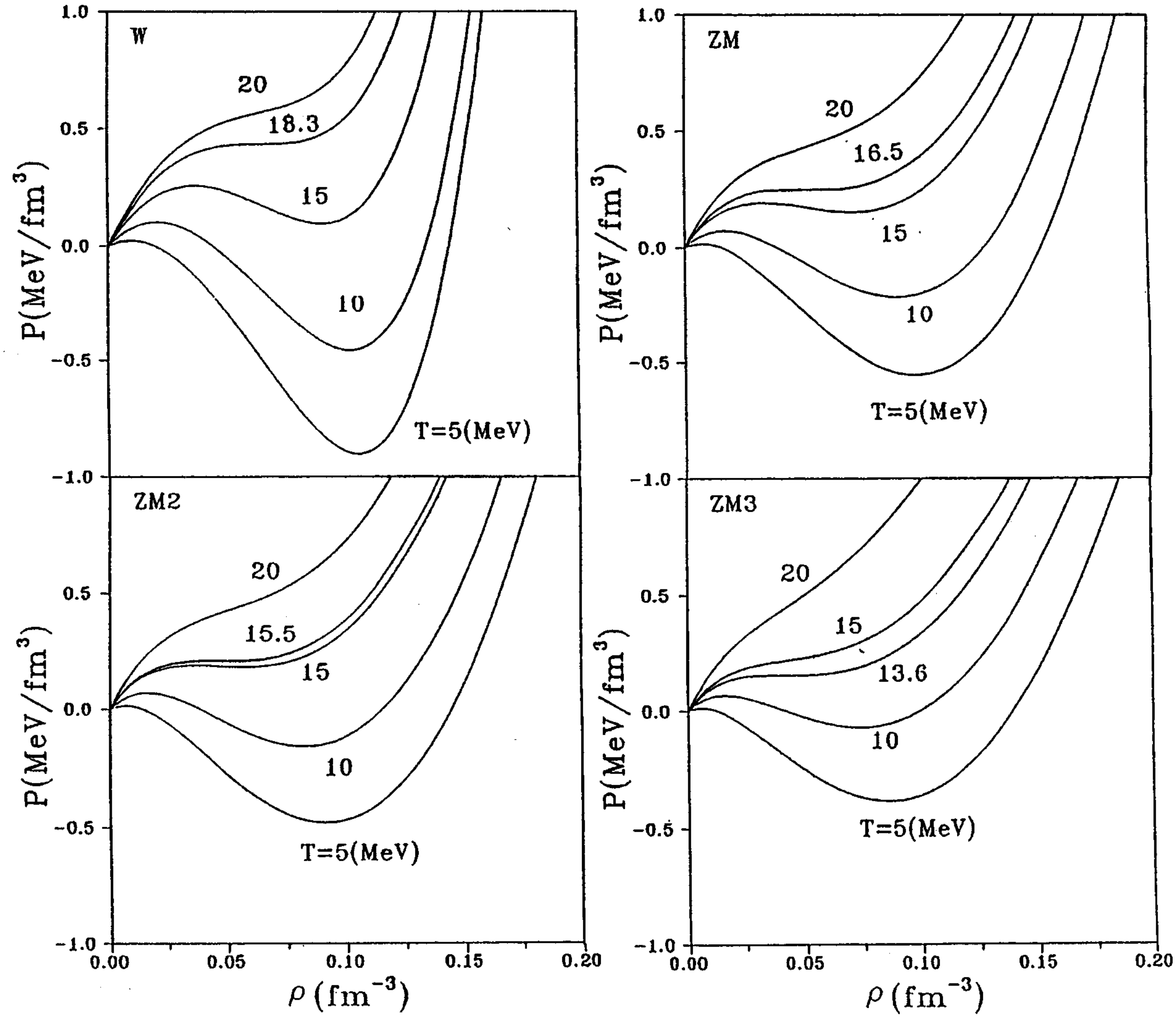

FIG. 4. Pressure as a function of baryon density at different temperatures for the Walecka model (W) and Zimanyi-Moszkowski models (ZM, ZM2, ZM3).

perature dependence. As the temperature is raised, M^* first increases and then decreases more slowly for the ZM models than for the Walecka model at $T=200$ MeV. Within the ZM models, this decrease is more pronounced in ZM3, but is even smaller compared to the Walecka model where the effective mass decreases very fast. In short, the effect of the temperature on the effective nucleon mass in the ZM models is not so pronounced as in the case of Walecka model, and can be seen only for densities below the normal density.

We present the energy per nucleon as a function of the density at various temperatures in Fig. 3. As the temperature increases, the nuclear matter becomes less bound and the the saturation curve around the equilibrium point in the ZM models is flatter than that in the Walecka model. This indicates that the nuclear matter EOS in the ZM models is softer compared to that obtained in the Walecka model, even at finite temperature. We can also conclude that the incompressibility of nuclear matter decreases when the temperature increases. This can be seen more clearly in Fig. 4 where we show the pressure-density isotherms of nuclear matter at different temperatures. Since the incompressibility K is related to $\partial p/\partial \rho$ (calculated at the equilibrium point where the pressure vanishes), we see directly that when the temperature increases, K decreases, and among the ZM models, the ZM3 model always gives the softest EOS for a fixed temperature.

The isotherms exhibit a typical van der Waals–like interaction where liquid and gaseous phases coexist. For very small temperatures the isotherms manifest the following behavior: for very low density the pressure increases with temperature as happens in an ideal gas, $p \sim \rho k_b T$. It decreases subsequently because of the attractive interaction of the sigma field, and finally increases as a consequence of the repulsion coming from the vector meson which dominates at high density. When the temperature increases, the term $\rho k_b T$ becomes more important and the local minimum in the pressure is less pronounced and disappears when the temperature is equal to the critical T_c. At this temperature, the unphysical region disappears and an inflection point appears in the isotherm, as we show in the Fig. 4 for each model. The $p-\rho$ isotherms in the ZM models have a shallower and flatter valley than the corresponding ones in the Walecka model, and this is more noticeable in the ZM3 model. In Table II we list the critical temperature T_c, density ρ_c, and pressure p_c given by the ZM and Walecka models. The ZM3 model pre-

TABLE II. Values for the critical temperature T_c and the effective mass M_c^* in MeV, critical density ρ_c in fm^{-3}, and pressure p_c in $\mathrm{MeV/fm}^3$ for the indicated models.

Models	T_c	ρ_c	p_c	M_c^*
Walecka	18.3	0.0650	0.4300	760
ZM	16.5	0.0698	0.2570	861
ZM2	15.5	0.0364	0.2106	881
ZM3	13.6	0.0354	0.1571	831

sents the lowest $T_c = 13.6$ MeV, density $\rho_c = 0.037\ \mathrm{fm}^{-3}$, and pressure $p_c = 0.157\ \mathrm{MeV\ fm}^{-3}$.

We present in Fig. 5 the phase diagram $T \times \rho$ of the models. The phase coexistence boundary is obtained when the liquid and gas phases have equal temperatures, chemical potentials, and pressures. Below the coexistence curve, the equilibrium state is a mixture of gas and liquid. This region is bigger in the Walecka model. In fact, if we include nonlinear terms in this model, this region becames smaller and the critical temperature goes down to $T_c = 14.2$ MeV [8]. The ZM3 model, where the non linearity of the coupling between the vector field to the nucleon is strongest, presents the smallest phase coexistence region compared to the other models.

As we have already pointed out, the nuclear matter incompressibility K decreases when the temperature increases. Therefore, we will have a temperature where the incompressibility K calculated at the equilibrium point vanishes. This temperature is known as the flash temperature $T = T_f$, $\partial p/\partial\rho|_{T_f} = p(\rho_f, T_f) = 0$. It represents the highest temperature at which a self-bound system can exist in hydrostatic equilibrium $(p=0)$. Above this temperature the warm nuclear matter is unbound and starts expanding. We present in Fig. 6 the pressure as a function of baryon density at the flash temperature for the models. This temperature is 14.1,

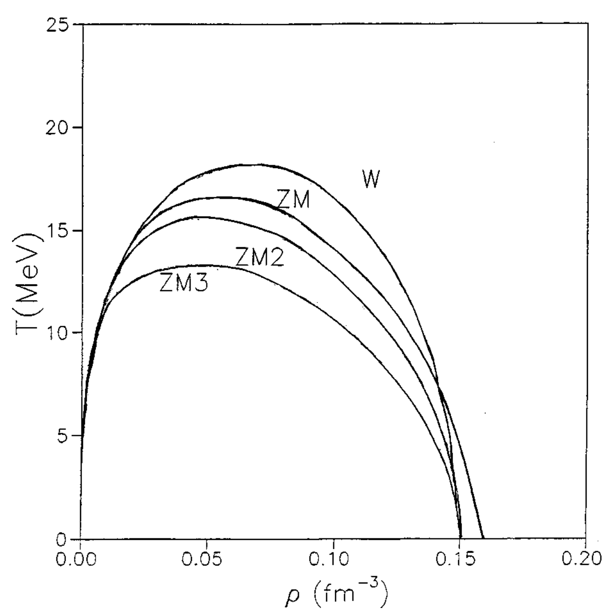

FIG. 5. Temperature as a function of the baryon density (phase diagram) for the Walecka model (W) and Zimanyi-Moszkowski models (ZM, ZM2, ZM3).

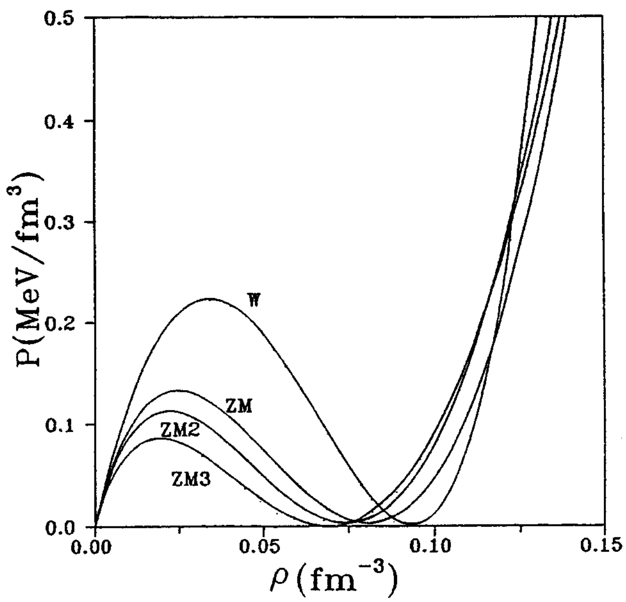

FIG. 6. Pressure as a function of baryon density at flash temperature (T_f) for the Walecka model (W) and Zimanyi-Moszkowski models (ZM, ZM2, ZM3).

12.9, 12.2, and 11.0 MeV for the Walecka, ZM, ZM2, and ZM3 models, respectively. Again, the ZM3 model has the smallest flash temperature. As expected, all of these temperatures are lower than the critical ones, because, as Fig. 4 shows, at the critical temperature the pressure is already positive and the system is expanding.

Finally, we present in Fig. 7 the entropy density as a function of the density at different temperatures. For high temperatures ($T = 200$ MeV), we see an increase in the entropy density with the density for all the models. This happens even at very low densities, and manifests what we have already pointed out when we discussed the behavior of the effective nucleon mass with the temperature at zero density. This decrease of M^* or increase of the entropy density with increasing temperature, which is more pronounced in the Walecka and ZM3 models, resembles a phase transition.

In summary, we have presented the thermodynamic properties of nuclear matter in three different versions of the ZM model. We have shown how the effective nucleon mass M^*, energy per nucleon, pressure, and entropy behave as a function of the density for different temperatures. As in the zero-temperature case, all the ZM models give a softer EOS of nuclear matter at finite temperature than the Walecka model. Among the three ZM models, ZM3 is the softest. Unlike the Walecka model the ZM models do not exhibit a phase transition for finite temperature at zero density. We studied the liquid-gas phase transition and found that these two new variants of the ZM model that incorporate a nonlinear new interaction between the nucleon and the vector meson present a much smaller phase coexistence region. The ZM3 model, in which the new interaction is stronger, presents the smallest phase coexistence region with the lowest critical temperature, density, and pressure. The incompressiblity decreases with increasing temperature, and vanishes when $T = T_{\mathrm{flash}}$. Again, the ZM3 model has the smallest flash temperature. Then we conclude that this new nonlinear interaction, which is so important to reproduce at zero-temperature

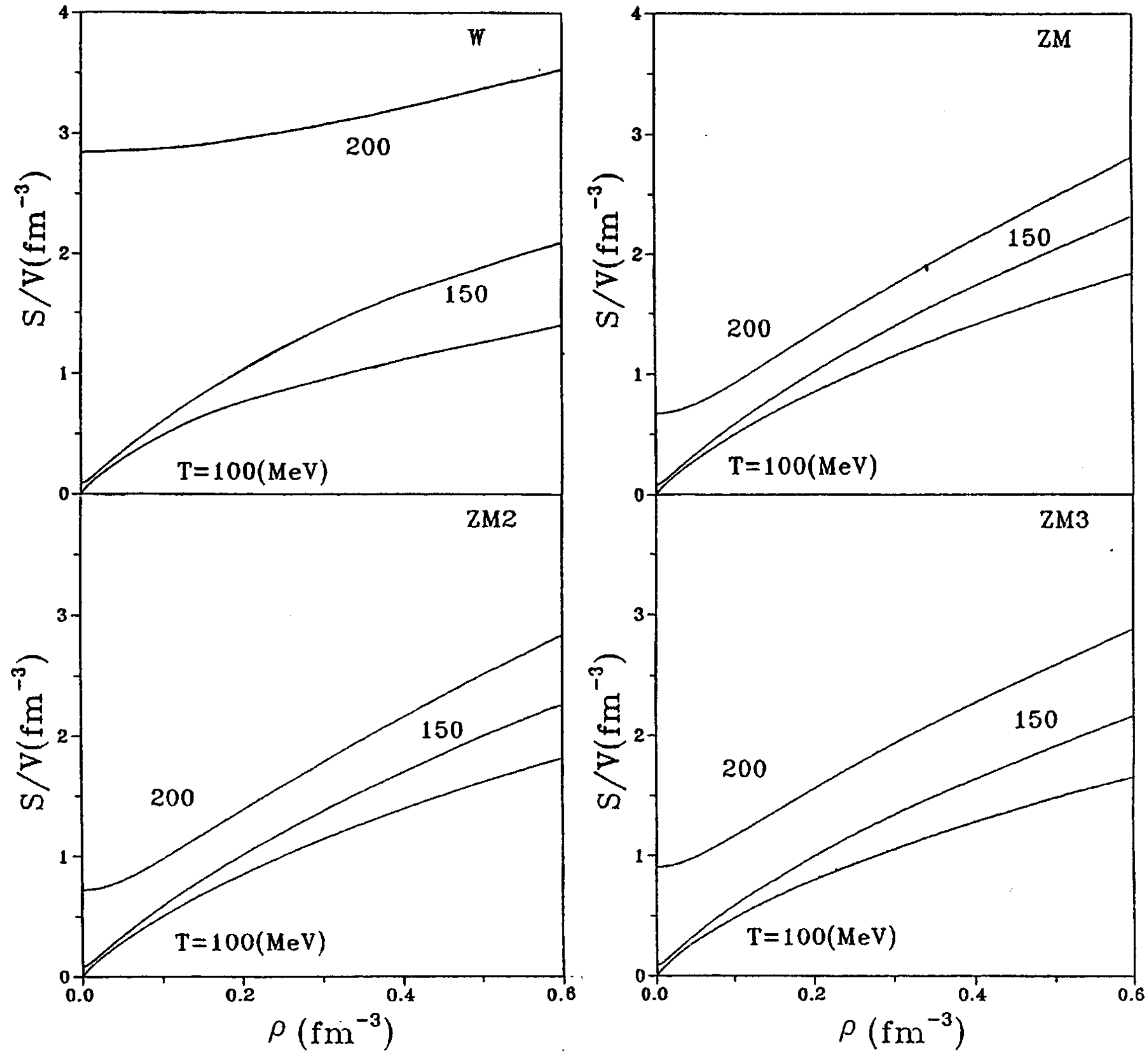

FIG. 7. Entropy density as a function of baryon density at different temperatures for the Walecka model (W) and Zimanyi-Moszkowski models (ZM, ZM2, ZM3).

nuclear matter and finite nuclei phenomenology [11,18], it is also very important at finite temperature. The experimental results investigating this warm and diluted matter produced in energetic heavy-ion collisions show a small liquid-gas phase region and a low critical temperature [4,1]. This suggests that the good results obtained in the ZM3 model at zero temperature remain even at finite temperature, and make the ZM3 model the most suitable of all models used.

ACKNOWLEDGMENTS

One of us (M.M.) would like to thank W. Melnitchouk for useful suggestions and a careful reading of the manuscript, the Maryland TQHN Group, supported by the U.S. Department of Energy, for their hospitality during his extended stay, and gratefully the Brazilian agency CAPES for the financial support which made this visit possible.

[1] T. Li *et al.*, Phys. Rev. C **49**, 1630 (1994) (they extracted a critical value for the slope parameter that converges for $T_c = 13.1 \pm 0.6$ MeV).

[2] J. E. Finn *et al.*, Phys. Rev. Lett. **49**, 1321 (1982); Phys. Lett. **118B**, 458 (1982); M. Mahi *et al.*, Phys. Rev. Lett. **60**, 1936 (1988).

[3] J. B. Elliot *et al.*, Phys. Rev. C **49**, 3185 (1994); M. L. Gilkes *et al.*, Phys. Rev. Lett. **73**, 1590 (1994).

[4] ALADIN/LAND Collaboration, J. Pochodzalla *et al.*, Phys. Rev. Lett. **75**, 1040 (1995).

[5] S. J. Lee and A. Z. Mekjian, Phys. Rev. C **56**, 2621 (1997).

[6] Granddon D. Yen, Mark I. Gorenstein, Walter Greiner, and Shin Nan Yang, Phys. Rev. C **56**, 2210 (1997).

[7] P. Papazoglou, J. Schaffner, S. Schramm, D. Zschiesche, H. Stocker, and W. Greiner, Phys. Rev. C **55**, 1499 (1997).

[8] R. J. Furnstahl and B. D. Serot, Phys. Rev. C **41**, 262 (1990).

[9] H. Muller and B. D. Serot, Phys. Rev. C **52**, 2072 (1995).

[10] J. Zimanyi and S. A. Moszkowski, Phys. Rev. C **42**, 1416 (1990).

[11] A. Delfino, C. T. Coelho, and M. Malheiro, Phys. Lett. B **345**,

361 (1995); Phys. Rev. C **51**, 2188 (1995).
[12] B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. **16**, (1986).
[13] A. Bhattacharyya and S. Raha, Phys. Rev. C **53**, 522 (1996).
[14] A. Delfino, M. Chiapparini, M. Malheiro, L. V. Belvedere, and A. O. Gattone, Z. Phys. A **355**, 145 (1996).
[15] Hua-lin Shi, Bao-qui Chen, and Zhong-yu Ma, Phys. Rev. C **52**, 144 (1995).
[16] G. Q. Li and Y. Z. Zhuo, Nucl. Phys. **A568**, 745 (1994).
[17] S. Gmuca, Nucl. Phys. **A547**, 447 (1992).
[18] M. Chiapparini, A. Delfino, M. Malheiro, and A. Gattone, Z. Phys. A **357**, 47 (1997).
[19] M. Malheiro, Mira Dey, A. Delfino, and Jishnu Dey, Phys. Rev. C **55**, 521 (1997).
[20] A. Delfino, J. Dey, M. Dey, and M. Malheiro, Phys. Lett. B **363**, 17 (1995); Z. Phys. C **71**, 507 (1996).
[21] Z. Qian, H. Song, and R. Su, Phys. Rev. C **48**, 154 (1993).
[22] J. Theis, H. Stocker, and J. Polony, Phys. Rev. D **28**, 2286 (1993).