Lawrence Berkeley National Laboratory

Recent Work

Title

FINITE TEMPERATURE METASTABLE MATTER

Permalink

<https://escholarship.org/uc/item/1hw8g5wp>

Author

Glendenning, N.K.

Publication Date

1986-10-01

Lawrence Berkeley Laboratory

UNIVERSITY OF CALIFORNIA

RECEIVED
LAWRENCE
BERKELEY LABORATORY

NOV 18 1986

LIBRARY AND
DOCUMENTS SECTION

Submitted to Physics Letters B

FINITE TEMPERATURE METASTABLE MATTER

N.K. Glendenning

October 1986

For Reference

Not to be taken from this room

DISCLAIMER

This document was prepared as an account of work sponsored by the United States Government. While this document is believed to contain correct information, neither the United States Government nor any agency thereof, nor the Regents of the University of California, nor any of their employees, makes any warranty, express or implied, or assumes any legal responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. Reference herein to any specific commercial product, process, or service by its trade name, trademark, manufacturer, or otherwise, does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof, or the Regents of the University of California. The views and opinions of authors expressed herein do not necessarily state or reflect those of the United States Government or any agency thereof or the Regents of the University of California.

Finite Temperature Metastable Matter†

Norman K. Glendenning

Nuclear Science Division
University of California
Lawrence Berkeley Laboratory
Berkeley, California 94720

October 7, 1986

Submitted to Physics Letters B.

†This work was supported by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

Finite Temperature Metastable Matter ¹

Norman K. Glendenning

Nuclear Science Division
University of California
Lawrence Berkeley Laboratory
Berkeley, California 94720

October 7, 1986

Abstract

Relativistic nuclear field theory possesses a solution corresponding to a hot metastable object of abnormal matter that cools by black-body radiation until a critical point is reached at which it becomes unstable, and would disassemble. If produced in a high energy nuclear collision, a large part of the initial energy of the object is emitted in black-body radiation, so that the nuclear fragments produced in its ultimate decay would account for only a fraction of the collision energy.

¹Work supported by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

Finite Temperature Metastable Matter

Norman K. Glendenning

October 7, 1986

1 Introduction

Relativistic nuclear field theory provides a good description of the bulk properties of nuclear matter as well as a large number of single-particle properties of finite nuclei [1]. With appropriate extensions the theory can be used to study matter away from the normal state, matter that is under extreme conditions of temperature or density, such as is expected to be produced in relativistic nuclear collisions, and as formed in the collapse of a star just prior to the bounce that produces the supernova, and as exists in the cores of the neutron stars into which the remaining matter of the star subsides. Being a non-linear theory in the field amplitudes, it has the possibility of possessing additional solutions to the normal state, that correspond to different field configurations [2]. In earlier work [3,4,5,6,7], a second solution was found at high temperature that is characterized by a low baryon effective mass and a high density of baryon-antibaryon pairs. Here we point out that under certain conditions there exists a temperature range over which matter in this abnormal phase can have zero pressure. In that case it is mechanically stable. Because of its internal properties, including its field configuration, there is a barrier against its decay to the normal state. An object made of such matter is therefore metastable. However since it is hot it will radiate. Because the baryons are far removed from their vacuum masses, there is a barrier against their emission. Instead, a fragment of such matter will cool by black-body radiation until it reaches a domain where its pressure becomes positive. It will then disassemble into baryons and perhaps lighter nuclei. If it is produced in a high energy nuclear collision at the top end of the range in temperature for which this metastable state can exist, most of the energy will be emitted in black-body radiation, and the nuclear fragments of its ultimate decay will account for only a fraction of the collision energy. It will

evolve in density during its radiation era from sub-nuclear to super-nuclear densities.

2 Theory

2.1 Lagrangian and Field Equations

As we shall discuss latter, abnormal solutions such as outlined above may exist for any theory having a scalar field coupled to baryons. However, the solution has special interest for the theory that can describe known nuclear properties. This is the scalar-vector theory [8,9,10], possibly augmented with ϕ^4 interactions [11]. For hot dense matter, the theory should be extended in several ways. First it is evident that in such matter, nucleons at the top of the Fermi sea are more energetic than other baryon species, including hyperons. So the excited nucleons, deltas and hyperons should also be incorporated. In hot matter, thermal mesons will also be produced. We made these generalizations in earlier studies. In non-strange matter the hyperons must appear as particle-antiparticle pairs. This would be a constraint on the solution in the case that the matter has a very short lifetime. A metastable object can develop a net strangeness with the decay of the associated kaon on the weak interaction time scale. This will lead to a still more stable configuration.

The extended Lagrangian is [3,4,5]

$$\mathcal{L} = \sum_B \bar{\psi}_B \left[i\gamma_\mu (\partial^\mu + ig_{\omega B} \omega^\mu) - (m_B - g_{\sigma B} \sigma) \right] \psi_B - \frac{1}{3} b m_n (g_\sigma \sigma)^3 - \frac{1}{4} c (g_\sigma \sigma)^4 + \mathcal{L}_\sigma^0 + \mathcal{L}_\omega^0 + \dots \quad (1)$$

Here ψ_B denotes a baryon spinor and the sum is over all of the baryon families, N, Δ , Λ , Σ , Ξ ... and their excited states until the solution converges over the range of densities considered. The $\mathcal{L}^0$ are the free Lagrangians of the mesons. The σ - and ω -mesons are Yukawa coupled to the baryon scalar density and vector current. The ρ -meson is coupled to the total isospin current, which however vanishes in symmetric matter. The other mesons are included as free thermal bosons, and their Lagrangians are represented by the ellipsis.

The field equations for static uniform matter in the mean field approximation are, for the mesons,

$$m_\sigma^2 \sigma = -b m_n (g_\sigma \sigma)^2 - c (g_\sigma \sigma)^3 + \sum_B g_{\sigma B} \langle \bar{\psi} \psi \rangle \quad (2)$$

$$m_\omega^2 \omega_0 = \sum_B g_{\omega B} \langle \psi^\dagger \psi \rangle \quad (3)$$

$$m_\omega^2 \omega_k = \sum_B g_{\omega B} \langle \bar{\psi} \gamma_k \psi \rangle = 0 \implies \omega_k \equiv 0 \quad (4)$$

and for the baryons,

$$\left[\gamma_\mu (p^\mu - g_{\omega B} \omega^\mu) - (m_B - g_{\sigma B} \sigma) \right] \psi_B = 0. \quad (5)$$

The baryon eigenvalues of momentum, p , for particle and antiparticle are,

$$\epsilon_B(p) = E_B(p) + g_{\omega B} \omega_0 \quad (6)$$

$$\bar{\epsilon}_B(p) = E_B(p) - g_{\omega B} \omega_0, \quad (7)$$

with

$$E_B(p) = \sqrt{p^2 + (m_B - g_{\sigma B} \sigma)^2}. \quad (8)$$

The source of attraction in the theory is the scalar field which reduces the masses of the baryons in the medium. We denote the effective mass by,

$$M_B = m_B - g_{\sigma B} \sigma. \quad (9)$$

In baryon matter, as contrasted with antibaryon matter, the vector field, eq. (3), is positive, being driven by the net baryon density. Therefor the vector field shifts the baryon eigenvalues, eq. (6) to higher energy, in baryon rich matter, and consequently is a source of repulsion. The effect on the antibaryons is opposite. When this shift is summed over the baryon spectrum, the result is a repulsive energy that is quadratic in the density.

2.2 Finite Temperature

For finite temperature the expectation values of the source currents appearing in the field equations are appropriate thermal averages over the Fermi-Dirac distribution functions for baryon and antibaryon [3,12],

$$n_B(p) = \frac{1}{1 + e^{(\epsilon_B(p) - \mu_B)/T}} \quad (10)$$

$$\bar{n}_B(p) = \frac{1}{1 + e^{(\bar{\epsilon}_B(p) + \mu_B)/T}}, \quad (11)$$

where, for antibaryons, $\bar{\mu}_B = -\mu_B$. The baryon chemical potential for species B ($= N, \Delta, \Lambda \dots$), is denoted by μ_B , and is,

$$\mu_B = \mu, \quad \text{for non - strange baryons} \quad (12)$$

$$\mu_B = g_{\omega B} \omega_0, \quad \text{for strange baryons} \quad (13)$$

Referring to eqs. (6,7,10,11), one sees that the latter condition is required to assure that the hyperons occur in particle-antiparticle pairs. The equations for σ and ω mesons are then,

$$m_\sigma^2 \sigma = -bm_n(g_\sigma \sigma)^2 - c(g_\sigma \sigma)^3 + \sum_B \frac{(2I_B + 1)(2J_B + 1)}{2\pi^2} g_{\sigma B} \int_0^\infty \frac{m_B - g_{\sigma B} \sigma}{\sqrt{p^2 + (m_B - g_{\sigma B} \sigma)^2}} (n_B(p) + \bar{n}_B(p)) p^2 dp \quad (14)$$

$$m_\omega^2 \omega_0 = \sum_B \frac{(2I_B + 1)(2J_B + 1)}{2\pi^2} g_{\omega B} \int_0^\infty (n_B(p) - \bar{n}_B(p)) p^2 dp. \quad (15)$$

Notice that σ occurs on the right side of both equations, since it is contained in the eigenvalue appearing in the Fermi-Dirac functions and also appears explicitly in the factor multiplying them. This is a highly non-linear equation, since if expanded in a power series in σ , it is of infinite order. The non-linearity arising from the potential terms with coefficients b and c are not essential to the phenomena that are discussed in this paper.

When the self-consistent solution for σ and μ are obtained at chosen temperature and baryon number density,

$$\rho = \sum_B \int_0^\infty (n_B(p) - \bar{n}_B(p)) p^2 dp, \quad (16)$$

the vector field, ω_0 , and the pressure and energy density can be calculated. The latter are given in this theory by the expressions

$$\begin{aligned}
p = & -\frac{1}{3}bm_n(g_\sigma\sigma)^3 - \frac{1}{4}c(g_\sigma\sigma)^4 - \frac{1}{2}m_\sigma^2\sigma^2 + \frac{1}{2}m_\omega^2\omega_0^2 \\
& + \frac{1}{3}\sum_B \frac{(2I_B+1)(2J_B+1)}{2\pi^2} \\
& \int_0^\infty \frac{p^2}{\sqrt{p^2+(m_B-g_{\sigma B}\sigma)^2}} (n_B(p) + \bar{n}_B(p)) p^2 dp \\
& + \frac{1}{3}\sum_M \frac{(2I_M+1)(2J_M+1)}{2\pi^2} \int_0^\infty \frac{p^2}{\sqrt{p^2+m_M^2}} \frac{p^2 dp}{e^{\sqrt{p^2+m_M^2}/T} - 1} \quad (17)
\end{aligned}$$

$$\begin{aligned}
\epsilon = & \frac{1}{3}bm_n(g_\sigma\sigma)^3 + \frac{1}{4}c(g_\sigma\sigma)^4 + \frac{1}{2}m_\sigma^2\sigma^2 + \frac{1}{2}m_\omega^2\omega_0^2 \\
& + \sum_B \frac{(2I_B+1)(2J_B+1)}{2\pi^2} \\
& \int_0^\infty \sqrt{p^2+(m_B-g_{\sigma B}\sigma)^2} (n_B(p) + \bar{n}_B(p)) p^2 dp \\
& + \sum_M \frac{(2I_M+1)(2J_M+1)}{2\pi^2} \int_0^\infty \sqrt{p^2+m_M^2} \frac{p^2 dp}{e^{\sqrt{p^2+m_M^2}/T} - 1} \quad (18)
\end{aligned}$$

In these equations, σ and ω_0 denote the mean values of the scalar meson, and the time-like component of the ω -meson. The space-like components vanish in isotropic matter. The sum on B is over the baryon species, and the one on M is over the thermal mesons, of which the pion triplet is most important.

We notice that if the sign of the chemical potential, μ , is everywhere changed, the sign of the baryon density and ω_0 is changed. However the scalar field equation remains unchanged as do p and ϵ . Therefor the solutions are reflection symmetric between positive and negative net baryon density (matter and anti-matter). This may be kept in mind when viewing graphs of properties as a function of density.

2.3 Parameters

The five parameters in the theory, g_σ/m_σ , g_ρ/m_ρ , g_ω/m_ω , b , c , are chosen so that the theory possesses the bulk properties of uniform symmetric matter, $B/A = 15.95$ MeV, saturation density $\rho = 0.145 fm^{-3}$, symmetry energy coefficient $a_{sym} = 36.8$ MeV [13], the compression modulus $K = 240$ MeV, and the nucleon effective mass at saturation, which we assume to be 0.8.

Fig. 1 Isotherms of energy density as a function of net baryon density. At low temperature, only the normal state exists, while at high temperature, only the abnormal. At intermediate temperatures, around $T = 150$ MeV, both solutions exist in a finite density interval.

The value of the compression modulus is consistent with the analysis of the giant monopole resonance [14], with the droplet model of atomic masses [13], and with the charge-distribution differences of heavy isotopes [15]. It is also consistent with known neutron star masses [16]. The corresponding coupling constants are $(g_\sigma/m_\sigma)^2 = 9.637 fm^2$, $(g_\omega/m_\omega)^2 = 4.5316 fm^2$, $(g_\rho/m_\rho)^2 = 6.4792 fm^2$, $b = 0.00895$, $c = 0.003689$. Of course nuclear matter properties do not determine the hyperon couplings. For simplicity we have assumed universal coupling. All well established baryons up to mass 1775 MeV were included in the calculation, and the mesons that were included in the thermal ensemble were similarly summed to convergence, although for the temperatures covered in this study, the pion is the most important.

3 High Temperature Properties

In the region in temperature and density above the gas-liquid phase transition, one might have expected that the energy density would be a monotonic increasing function of baryon density. However, as noticed previously [3], there is a region in which three solutions exist at the same density. On the isotherms shown in Fig. 1, this can be seen for $T = 150$ and 160 MeV. Two of the three solutions are, in the case of uniform matter, local minima in the thermodynamic potential regarded as a function of σ and the other is a maximum and therefore an unstable solution (the middle one). We have not

Fig. 2 Same as Fig. 1, but a finer mesh on temperature, showing the details of the onset of the abnormal phase.

examined stability with respect to space-varying fields. However, we cannot find any source of attraction in the theory arising from a spatial variation that would more than compensate for the corresponding kinetic energy, so we believe that the minima found for uniform matter are true minima of the thermodynamic potential.

The onset of the region of two phases is shown on a finer grid of isotherms in Fig. 2. In particular one sees that below about $T = 146$ MeV, only the normal phase exists. At $T = 148$ MeV, a well developed second order phase transition exists. The lower energy branch is the normal state of hot nuclear matter. Above about $T = 148$ MeV, this phase is confined to the lower density region. The upper branch is an abnormal state, characterized by low nucleon effective mass and an abundance of baryon-antibaryon pairs, both strange and non-strange. The branch that connects them is unstable, as it must be since all solutions correspond to the stationary points on the surface of the thermodynamic potential, so that stable and unstable solutions must alternate. Referring again to Fig. 1, by $T = 150$ MeV, the abnormal state has become disjoint from the other solutions, and with increasing temperature the density domain of the normal solution shrinks until it vanishes around $T = 160$ MeV. Above this temperature only the abnormal state exists for any baryon density. The high energy density in the abnormal state at zero, or small (net) baryon density, corresponds to the abundance of pairs.

We emphasize here that the precise domain in temperature and density

where these phenomena occur depends sensitively on the value chosen for the nucleon effective mass at saturation. This is not a well known quantity. The value chosen here is $m^*/m = 0.8$. If smaller, the phenomena move to a lower domain. The domain of validity of the theory is not well known, but it is terminated at high density or temperature by the finite size of baryons and their quark structure. To be a realizable state of nature, the temperature and density of the abnormal phase must not be too high.

3.1 Origin of the Abnormal State

The abnormal state discussed here is a finite temperature one, unlike the one discussed by Lee and Wick [17]. It is the finite temperature that drives the state, in the following way. At finite temperatures, baryon-antibaryon pairs can be created, because they have the quantum numbers of the vacuum. Baryon and antibaryon act differently on the scalar and vector fields. Being scalar, the scalar field is driven equally by baryon and antibaryon, eq. (14). The pairs therefore increase the scalar field strength, reducing their effective mass, thereby making it energetically possible to accommodate even more pairs. Thus an instability can arise, leading to a second order phase transition. The vector field, on the other hand, being vector, *can* distinguish baryon and antibaryon, and the pair has null effect on the vector field, eq. (15). These arguments, as was earlier emphasized, are not particular to the precise form of the Lagrangian, so long as there is a scalar field coupled to the baryons [5]. The principal characteristics of the abnormal state, are the large scalar field, leading to low effective baryon masses, the high abundance of baryon pairs, and the corresponding high entropy per net baryon.

3.2 Metastable Bodies

In a narrow range of temperatures, $147 \text{ MeV} < T < 152 \text{ MeV}$, the pressure becomes zero over a segment of the abnormal branch. This is shown for $T = 148 \text{ MeV}$ in Fig 3. The point of zero pressure on the abnormal branch is mechanically stable. Clearly the source of binding energy is the high abundance of baryon-antibaryon pairs that are characteristic of the abnormal state. They reduce the effective masses, while having null effect on the repulsive vector field.

Let us suppose that a piece of matter has been produced in this state, in a high energy collision between heavy nuclei. The energy required to produce the abnormal state at $T = 151 \text{ MeV}$ in the zero pressure configuration

Fig. 3 Pressure as a function of density on the isotherm $T = 148$ MeV., showing details of the normal, abnormal and unstable solution. Note that on the unstable solution the pressure is zero at a density of about 0.6 fm^{-3} .

is 22 GeV per nucleon in the center of mass system. Having zero pressure, it is mechanically stable. The surface tension and Coulomb force will actually displace it slightly to a position of positive pressure, but for our purpose, we can ignore this. It is metastable because its field configuration is very different from the normal state of matter, that is to say, the scalar field is very large. The baryon masses are far displaced from their vacuum values or their values in the normal state of matter. Moreover, there is a high abundance of low mass baryon-antibaryon pairs. The entropy is correspondingly very high. All of these factors introduce a barrier to its decay to the normal state. However, since the object is hot, $T \approx 150$ MeV, and since the baryons are far removed from their vacuum masses, it will radiate bosons in a spectrum characteristic of its temperature. These will include photons, lepton pairs, and pions. Since it is dense, it is plausible that the strong interactions will maintain a thermal distribution of bosons as it cools. Its spectrum will therefore be that of a very hot black-body. In Fig. 4 the energy per net nucleon as a function of temperature in this zero pressure configuration of the abnormal state is shown. The specific heat is obviously very high, requiring 6 GeV in radiation to cool it from 151 to 150 MeV. About two thirds of the energy of the body is emitted in black-body radiation, as it cools to $T = 147$ MeV. In the vicinity of that temperature, the abnormal state ceases to have negative pressure, and somewhat below that temperature, the distinction between normal and abnormal phase ceases (see Fig. 2). Presumably the body will commence to expand and disassemble into baryons and light particle fragments. However, because of its high entropy in the abnormal state,

Fig. 4 Center of mass energy per nucleon as a function of temperature for zero-pressure abnormal matter.

XBL 8610-3762

the expansion cannot be a smooth hydrodynamic one. The matter must pass through a shock discontinuity in order to conserve or increase entropy as it passes to the normal state. The energy distribution of the produced particles will not be that expected of a high energy collision because of the large fraction of the energy that was carried off in black-body radiation. Only one third of the original energy is available at the disassembly stage.

4 Summary

Because of their non-linearity, the field equations of nuclear field theory admit of additional solutions besides the normal state of matter. One of these is an abnormal state, characterized by low baryon effective masses, abundant baryon-antibaryon pairs, and high entropy. Over a narrow range in temperature, matter can exist in the abnormal state at zero pressure. This is a hot metastable state, for which there is a barrier against decay, because the field configuration is different than the normal state, the baryon masses are far removed from their vacuum masses, there is an abundance of pairs also far removed from their vacuum masses, and a correspondingly high entropy. The abundance of baryon-antibaryon pairs is the glue that holds this matter together. The signals associated with this novel state are quite unusual. A fragment of such matter will cool by emitting a spectrum of black-body radiation, consisting principally of photons, lepton pairs and pions, rather than by baryon emission, because they are far removed from

their vacuum masses. If produced at the upper end of its temperature range, a large fraction of the original energy is radiated in this way. The baryons and light elements produced in the eventual decay will account for only a fraction of the original energy.

Two aspects of the metastable state that we have not yet investigated are its lifetime, and the increase in stability that it would acquire if its lifetime were long on the scale of the weak interactions, through the development of a net strangeness. For a short-lived object produced in the collision of ordinary nuclei, the hyperons must appear in pairs. For a long lived object, some of the conserved baryon charge can be carried by hyperons, which is energetically favored when the chemical potential exceeds their effective masses. If sufficiently long-lived, an additional signal would be the production of hyperons in the ultimate decay of the metastable fragment.

Concerning the prospects for the existence in nature of such a novel state as discussed here, we note that nuclear field theory is an effective one, whose range of validity is terminated at high density by the finite size of the nucleons and their underlying quark structure. Its precise range of validity is unknown. Secondly, the precise location in temperature and density of the metastable state depends sensitively on the nucleon effective mass, which is not well known. For the value of the effective mass, $m^*/m = 0.8$, chosen here, the temperature and density of the metastable state is on the marginally high side, but may still lie below the phase transition to the quark-gluon plasma. If the effective mass is lower than the above value, the abnormal phase would lie very comfortably in the lower temperature-density domain. In any case, the theory, regarded as a mathematical model, does possess such a solution, and it would be a very novel state of matter were it to be found.

Acknowledgements: I appreciate a critical reading of the manuscript by F. Klinkhammer and W. Swiatecki. This work was supported by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

References

- [1] P.-G. Reinhard, M. Rufa, J. Marhun, W. Greiner and J. Freidrich, Z. Phys. **A323** (1986) 13.
- [2] B. Banerjee N. K. Glendenning and M. Gyulassy, Nuc. Phys. **A361** (1981) 326.
- [3] S. I. A. Garpman, N. K. Glendenning and Y. J. Karant, Nuc. Phys. **A322** (1979) 382.
- [4] N. K. Glendenning, *Phase Transitions in Nuclear Matter* in 7th High Energy Heavy Ion Study, 1984, ed. R. Bock, H. H. Gutbrod and R. Stock (GSI report-85-10).
- [5] N. K. Glendenning, Phys. Lett. **144B** (1984) 158.
- [6] J. Theis, G. Grabner, G. Buchwald, J. Maruhn, W. Greiner, H. Stocker, J. Polonyi, Phys. Rev. **D28** (1983) 2286.
- [7] T. Nakai and S. Takagi, Prog. Theor. Phys. **71** (1984) 1118.
- [8] M. H. Johnson and E. Teller, Phys. Rev. **98** (1955) 783;
- [9] H. P. Duerr, Phys. Rev. **103** (1956) 469;
- [10] J. D. Walecka, Ann. of Phys. **83** (1974) 491
- [11] J. Boguta and A. R. Bodmer, Nucl. Phys. **A292** (1977) 413.
- [12] J. D. Walecka, Phys. Lett. **59B** (1975) 109.
- [13] W. D. Myers, *Droplet Model of Atomic Nuclei* (New York: McGraw Hill, 1977); W. D. Myers and W. Swiatecki, Ann. of Phys. **55** (1969) 395; W. D. Myers and K-H. Schmidt, Nucl. Phys. **A410** (1983) 61.
- [14] J. P. Blaizot, D. Gogny and B. Grammaticos, Nucl. Phys. **A265** (1976) 315; J. P. Blaizot, Phys Rep. **64** (1980) 171; J. Treiner, H. Krevine, O. Bohigas and J. Martorell, Nucl. Phys. **A317** (1981) 253.
- [15] G. Co and J. Speth, Phys. Rev. Lett. **57** (1986) 547.
- [16] N. K. Glendenning, Phys. Rev. Lett. **57** (1986) 1120.
- [17] T. D. Lee and G. C. Wick, Phys. Rev. **D9** (1974) 2291.

This report was done with support from the Department of Energy. Any conclusions or opinions expressed in this report represent solely those of the author(s) and not necessarily those of The Regents of the University of California, the Lawrence Berkeley Laboratory or the Department of Energy.

Reference to a company or product name does not imply approval or recommendation of the product by the University of California or the U.S. Department of Energy to the exclusion of others that may be suitable.

*LAWRENCE BERKELEY LABORATORY
TECHNICAL INFORMATION DEPARTMENT
UNIVERSITY OF CALIFORNIA
BERKELEY, CALIFORNIA 94720*