Condensation States And Landscaping With The Theory Of Abstraction

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Abstract:

The Abstraction theory is applied in landscaping. A collection of objects may be made to be vast or meager depending upon the scale of observations. This idea may be developed to unite the worlds of the great vastness of the universe and the minuteness of the sub-atomic realm. Keeping constant a scaling ratio for both worlds, these may actually be converted into two self-same representatives with respect to scaling. The Laws of Physical Transactions are made use of to study Bose-Einstein condensation. As the packing density of concerned constituents increase to a certain critical value, there may be evolution of energy from the system.

Introduction:

Be it the large vastness of the universe or the delicate smallness of the sub-atomic world, by choosing a suitable constant scaling ratio for both, we may obtain their representations. These representations following a certain constant scaling ratio, will be self-same. In previous papers on the subject, I have mentioned the chaotic behavior in the quantum world. Choosing suitable scaling ratios, we may turn the universe itself into such a chaotic quantum system, having its own necessary quantum states and trajectory behaviour. In that case, the study of the universe reduces to the study of some sort of a quantum chaotic system. On the other hand, choosing some other necessary scaling ratios, the atomic and the sub-atomic realm may be extended to become the universe itself, complete with its own macroscopic trajectory behaviour. Instead of formulating different ways of looking at worlds of different sizes, if we adjust the way of viewing i.e., the scaling ratio in such a fashion that the representations of the world merge, we will be looking at representative worlds of study which are practically self-same. The Laws of Physical Transactions formulated in previous papers of the subject may then be applied in order to study such self-same representations of the worlds of various scales. Unification of the ways of studying at different ranges of scaling may thus be achieved by suitable landscaping (adjusting different scales to a suitable scaling-ratio, in order to make all the scales of study similar in size). Further, a similar approach may be applied to study the Bose-Einstein Condensation. A certain critical packing density of the constituents of each world of a certain landscape must ensure a condensation of similar sort. The quantum states (or some similar states) of each such landscape will merge and give spikes for that critical scaling ratio in their respective representations.

The quantum chaotic behavior may be of interest to study if we are to learn about the universe as a whole. The astronomically large distances separating clusters in the universe supports a study of such sorts. Quantum chaotic behaviour, on the other hand will give rise to something similar to the Bose-Einstein condensation at some critical packing density. The study of such condensation states too will be of interest here.

Scaling The Universe:

Looking at a large enough part of the universe, we may draw an analogy to a system of scattered particles in motion or rest relative to each other. These particles may or may not be similar to each other, if we look at a given locality. Our idea, however, is that we can always represent even the whole of the universe on a piece of paper of our desired size. We can very well do the same with localities of sub-atomic sizes.

We may represent both the worlds, viz. the microscopic and the macroscopic, within any desired standard size. Theoretically, we are only to diminish the snaps of the universe and magnify the snaps of the microscopic world in order to put both into representations of a definite scaling-size. Looking at such a representation of the macroscopic world (due to the large number of constituents and the large distances separating them involved) we will find it to be a complex mixture of various kinds of particles. On the other hand, looking at such a representation of the microscopic world, (due to the small distances separating the constituents) it will be like the actual universe itself, with various types of constituent parts involved. Such a representation of the microscopic and the macroscopic worlds will bring out hidden properties and behaviours of both worlds, as well as providing for a similar basis of studying them both.

Let us consider a given representation with fractal dimension D_F. The fractal dimension is purely geometrical, i.e., it only depends on the shape of the representation. A suitable probability measure $d\mu$, according to the particular phenomenon considered is assigned to the given representation. A coarse grained probability density, as the mass of the hypercube Λ_i of size l is defined as,

$$P_i(l) = \int_{\Lambda_i} d\mu\,(x) \qquad \dots (1);$$

where $i = 1,2,3, \dots, N(l)$.

The information dimension D_I is such that,

$$\sum_{i=1}^{N(l)} P_i \ln(P_i) \simeq D_I \ln(l) \qquad \dots (2);$$

where $D_I \leq D_F$.

The number of boxes containing the dominant contributions to the total mass and thus relevant part of the information, is,

$$N_R(l) \propto l^{-D_I} \qquad \dots (3).$$

For each box Λ_i, $D_I = D_F$ for a uniform distribution. When $D_I < D_F$, the measure itself may be called fractal since it is singular with respect to the uniform distribution,

$$P^* = \frac{1}{N(l)} \propto l^{D_F}$$

for each box Λ_i.Thus, $\frac{P_i}{P_i^*}$ can diverge in the limit of vanishing l.

Simulations of the mass-moment scaling yields,

$$\langle P_i(l)^q \rangle \equiv \sum_{i=1}^{N(l)} P_i(l)^{q+1} \propto l^{q.d_{q+1}} \qquad \dots (4).$$

The d_q are the Renyi dimensions which generalize the information dimension $D_I = d_1$as well as the fractal dimension $D_F = d_0$. If the d_q's are not constant, anomalous scaling is to be employed and, as the order q varies, the amount of the difference $D_q - D_F$ gives a first rough measure of the heterogeneity of the probability distribution.

The moment generic observables A computed on scale l is such that,

$$\langle A(l)^q \rangle \propto l^{g(q)}$$

Anomalous scaling, i.e., a non-linear shape of the function $g(q)$ is the more common situation, where one does not require unnecessarily to consider only a finite number of scaling components. In some cases one may observe strong time variations in the degree of chaoticity. This intermittency phenomenon involves an anomalous scaling with respect to time-dialations identifying the parameter e^{-t} wit the parameter l used in spatial dialations. A measure of the degree of intermittency requires the introduction of infinite sets of exponents which are analogous to the Renyi dimensions and can be related to a multifractal structure given by the dynamical system in the functional trajectory space.

The Grassberger-Procaccia correlation dimension v is defined by considering the scaling of the correlation integral,

$$C(l) = \lim_{M \to \infty} \frac{1}{M^2} \sum_i \sum_{j \neq i} \theta\left(l - |x_i - x_j|\right);$$

where θ is the Heaviside step function and $C(l)$ is the percentage of pairs (x_i, x_j) with distance $|x_i - x_j| \leq l$.

In the limit $l \to 0$,

$$C(l) \propto l^v.$$

In general,

$$v \leq D_F.$$

v is a more relevant scaling index than D_F since it is related to the point probability distribution on the attractor, while D_F cannot take into account an eventual homogeneity in the visit frequencies.

Let us define the number of points in an F-dimensional spherical representation of the world, with radius l and centre at x_i as,

$$n_i(l) = \lim_{M\to\infty} \frac{1}{M-1} \sum_{j\neq i} \theta\,(l - |x_i - x_j|) \qquad \ldots(5).$$

We must introduce a whole set of generalized scaling exponents

$$\langle n(l)^q \rangle = \lim_{M\to\infty} \frac{1}{M} \sum_{i=1}^{M} n_i(l)^q \propto l^{\emptyset(q)} \qquad \ldots(6);$$

where $\emptyset(1) = v$.

Considering a uniform partition of phase space into boxes of size l it is convenient to introduce the probability $P_K(l)$ that a point x_i falls into the K^{th} box. In this case, the moments of P_K can be estimated by summing up the boxes,

$$\langle p(l)^q \rangle = \sum_{K=1}^{N(l)} P_K(l)^{q+1} \propto l^{q.d_{q+1}}.$$

A moment of reflection shows

$$\emptyset\,(q)/q = d_{q+1}$$

because of the ergodicity $n_i(l) \sim P_K(l)$, if x_i belongs to the K^{th} box and since one can use either an 'ensemble' average (weighted sum over the boxes) or a 'temporal' average (sum of the time evolution $x(l)$).

The fractal dimension for $q = -1$ is,

$$D_F = d_0 = -\emptyset(-1) \qquad \ldots(7)$$

while the correlation dimension is,

$$v = d_2 = \emptyset(1) \qquad \ldots(8)$$

Statistical laws at small scales have to depend not only on the average energy dissipation density $\bar{\varepsilon}$ but also on the fluctuations of energy dissipation density $\varepsilon(x)$.

According to the Theory of Physical Abstraction, each point x should have the same singularity structure,

$$\Delta V_x(r) \propto r^h, h = \frac{1}{3} \qquad \ldots(9)$$

In other words $\varepsilon(x)$ tends to be smoothly distributed in a region of R^3. The eddy turn-over time and the kinetic energy per unit mass at scale r are defined as,

$$t(r) \sim \frac{r}{\Delta V(r)} \qquad \ldots(10)$$

and $$E(r) \sim \Delta V(r)^2 \quad ...(11)$$

the transfer rate of energy per unit mass from the eddy at scale r to smaller eddies is then given by

$$\tilde{\varepsilon}(r) = \frac{E(r)}{t(r)} \sim \frac{\Delta V(r)^3}{r} \quad ...(12)$$

Since $\varepsilon_x(r) = \left(\frac{1}{r^3}\right) \int_{\Lambda_x(r)} \varepsilon(y) d^3y$, [$\Lambda_x(r)$ is a cube of edge r around x] we have,

$$\int_{\Lambda_x(r)} \varepsilon(y) d^3y \sim r^3 \quad ...(13)$$

$r \to 0$ means r in the initial range and the regions containing a large part of $\varepsilon(x)$ are a physical approximation of a fractal structure. In this β −model approach,

$$\int_{\Lambda_x(r)} \varepsilon(y) d^3y \propto \begin{cases} r^{D_F} & if\ x \in S \\ 0 & if\ x \notin S \end{cases}$$

in an equivalent way

$$\Delta V_x(r) \propto \begin{cases} r^h & if\ x \in S \\ 0 & if\ x \notin S; \end{cases}$$

where $h = (D_F - 2)/3$

At scale r,there is only a fraction,

$$r^{3-D_F} \propto \frac{r^{-D_F}}{r^{-3}}$$

occupied by active eddies.

The transfer energy from the eddy at scale l_n (active eddy) to the scale l_{n+1} is,

$$\varepsilon_n \propto \frac{v_n^3}{l_n}$$

Since, the energy transfer rate is constant in the cascade process, for $\beta = 2^{D_F-3}$, we have,

$$\varepsilon_n = \beta \varepsilon_{n+1}, \frac{v_n^3}{l_n} = \beta \frac{v_{n+1}^3}{l_{n+1}} \quad ...(14)$$

Iterating, we have,

$$v_n \propto l_n^{1/3} (l_n/l_0)^{(D_F-3)/3} \quad ...(15)$$

Each eddy at scale l_n is divided into eddies of scale l_{n+1} in such a way that the energy transfer for a fraction β of eddies increases by a factor $\frac{1}{\beta}$, while it becomes zero for the other ones.

In order to generalize the β-model, we have at scale l_n, N_n active eddies. Each eddy $l_n(k)$ generates active eddies covering a fraction of volume $\beta_{n+1}(k)$. k labels the mother-eddy and $k = 1, \dots, N_n$.

Since the rate of energy transfer is constant among mother-eddies and their effects, we have,

$$\frac{\nu_n(k)^3}{l_n} = \beta_{n+1}(k)\frac{\nu_{n+1}(k)^3}{l_{n+1}} \quad \dots (16)$$

The iteration of ν_n gives an eddy generated by a particular history of fragmentations $[\beta_1, \dots, \beta_n]$, such that,

$$\nu_n \propto l_n^{1/3}\left(\prod_{i=1}^{n}\beta_i\right)^{-1/3} \quad \dots (17)$$

The fraction of volume occupied by an eddy generated by $[\beta_1, \dots, \beta_n]$ is $\prod_{i=1}^{n}\beta_i$, such that,

$$\langle|\Delta V(l_n)|^P\rangle \propto l_n^{P/3}\int\prod_{i=1}^{n}d\beta_i\ \beta_i^{(1-P/3)}P(\beta_1, \dots, \beta_n) \quad \dots (18)$$

With no correlation among different steps of the fragmentation, i.e., with $P(\beta_1, \dots, \beta_n) = \prod_{i=1}^{n}P(\beta_i)$, the exponent concerned,

$$\zeta_P = \frac{P}{3} - \ln_2\{\beta^{(1-P/3)}\} \quad \dots (19)$$

Let us now consider a given representation of the universe. Let the packing density of the constituents be $\emptyset$. This packing density function $\emptyset$ will affect any given constituent point inside it in accordance with the Laws of Physical Transactions. The given constituent point concerned will in turn affect $\emptyset$ while interacting. For a given critical state of study of the total effects, we therefore are going to have a shear stress $\emptyset$ and a mean effective stress f. The critical state line is the loci of critical state conditions in the $\varepsilon - f - \emptyset$ space. Its projection on the $f - \emptyset$ space defines a strength parameter,

$$M = \frac{\emptyset}{f} = \frac{6\sin\emptyset}{3-\sin\emptyset} \quad \dots (20).$$

The second equality applies to axis-symmetric, axial compression and it is a function of the constant volume critical state packing density function $\emptyset$.

The small-strain stiffness of a given representation is measured by imposing a smaller strain than the elastic threshold strain concerned. In this range, deformations localize at inter-point contacts and the granular skeleton deforms at constant fabric of spacetime. The nonlinear load-deformation response determines the stress-dependent shear wave velocity

$$V_S \propto \left(\frac{f-\emptyset}{\varepsilon}\right)^{\beta} \quad \dots (21)$$

Inside a given cluster, we may have various growth patterns. The growth may occur mainly at an active zone on the surface of the cluster. For a one-dimensional interface, a fluctuation-dissipation theorem exists, leading to an exact dynamic exponent $z = \frac{3}{2}$. This is in excellent agreement with numerical simulations of ballistic aggregation and Eden clusters. For two-dimensional interfaces, $z \sim 1.5$.

The interface profile is described by a height $h(x,t)$. The simplest nonlinear Langevin equation for a local growth of the profile is,

$$\frac{\partial h}{\partial t} = m\Delta^2 h + \frac{\lambda}{2}(\nabla h)^2 + \eta(x,t) \qquad \ldots (22)$$

The first term on the right-hand side describes relaxation of the interface by a tension term m. The second term is the lowest -order nonlinear term that can appear in the interface growth equation. Higher-order terms may also be present, but they are irrelevant and will not modify the scaling properties concerned. The noise $\eta(x,t)$ has a Gaussian distribution with $\langle \eta(x,t) \rangle = 0$ and

$$\langle \eta(x,t)\eta(x',t') \rangle = 2D\delta^d(x-x')\delta(t-t').$$

There is also a velocity term, but it is removed by choice of an appropriate moving coordinate system. Equation (22) is invariant under translation $h \to h + constant$, and obeys the infinitesimal reparametrization,

$$h \to h + b.x, x \to x + \lambda bt,$$

which describes the tilting of the interface by a small angle. When a given constituent point is added, the increment projected along the h-axis is,

$$\delta h = m\sqrt{[1 + (\nabla h)^2]} \simeq m + (m/2)(\nabla h)^2 + \cdots$$

Following the transformation $W(x,t) = e^{[(\lambda/2m)h(x,t)]}$, we have

$$\frac{\partial W}{\partial t} = m\nabla^2 W + \left(\frac{\lambda}{2m}\right)\eta(x,t)W \qquad \ldots (23)$$

which is a diffusion equation in a time-dependent random potential. $W(x,t)$ is the sum of Boltzmann weights for all static configurations of a flow in a $(d+1)$-dimensional space from $(0,0)$ to (x,t). The noise term describes a quenched random potential $(\lambda/2m)\eta(x,t)$exerted by the environment. The second transformation, $m = -\nabla h$ results in

$$\frac{\partial m}{\partial t} + \lambda m\,.\,\nabla m = m\nabla^2 m - \nabla\eta(x,t) \qquad \ldots (24),$$

which is the Burger's equation for a vorticity-free velocity field for $\lambda = 1$. In the Burger's equation, further evolution of the pattern proceeds through the larger parabolas growing at the expense of the smaller ones, and parallels the evolution of shock waves.

If the initial profile is $h(x,0) = h_0(x)$, its evolution is given by,

$$h(x,t) = \frac{2m}{\lambda}\ln\left\{\int_{-\infty}^{\infty} \frac{d^d\xi}{(4\pi mt)^{\frac{d}{2}}} e^{\left[-\frac{(x-\xi)^2}{2mt} + \frac{\lambda}{2m}h_0(\xi)\right]}\right\} \qquad \ldots (25)$$

Let a given representation have bonds within itself, occupied by a resistance generated inside it due to its packing density $\emptyset$, with probability p. Let it have a support towards the concerned flow with probability $1-p$. In such a representation, we have,

$$\langle \sum_k \varepsilon_k^n \rangle_{\xi,L} \propto L^{-x_n}, L \to \infty \quad \dots (26);$$

where $\langle \sum_k \varepsilon_k^n \rangle$ refers to the average over the sample realizations, L is the system-size $L \lesssim \xi$ and $\xi \propto (p - p_c)$ is the correlation length. ε_k is the energy dissipated in the branch k.

For a finite size scaling behaviour,

$$p\left(\sum_k \varepsilon_k^0, \sum_k \varepsilon_k^1, \dots, \xi, L\right) = \lambda^{x_0}\lambda^{x_1} \dots p\left(\sum_k \varepsilon_k^0/\lambda^{-x_0}, \frac{\sum_k \varepsilon_k^1}{\lambda^{-x_1}}, \dots, \xi/\lambda, L/\lambda\right),$$

(λ is the rescaling parameter) equation (26) implies,

$$\langle \left(\sum_k \varepsilon_k^n\right)^m \rangle_{\xi,L} \propto L^{-m x_n} \quad \dots (27)$$

In disordered representations, the fluctuations of the free energy among different replicas may be regarded as the analogue of the temporal intermittency in a chaotic signal. Considering a spin-model of the D-dimensions, the Hamiltonian,

$$H[\{J_{ij}\}] = -\sum_{(i,j)} J_{ij}\sigma_i\sigma_j,$$

where $\sigma_i = \pm 1$ is the of the spin on the site i and the coupling J_{ij} is an independent random variable distributed according to a probability distribution $p(J_{ij})$. Given a coupling realization $\{J_{ij}\}$, the partition function of an N spin system is the trace of the Boltzmann weight $e^{(-\beta H_N)}$,

$$Z_N(\beta, \{J_{ij}\}) = \sum_{\{\sigma_i\}} e^{\{-\beta H_N[\{J_{ij}\}]\}}$$

The free energy per spin in the limit $N \to \infty$ is,

$$F(\beta) = \lim_{N\to\infty} -\frac{1}{N\beta}\langle \ln Z_n \rangle = \lim_{N\to\infty} -\frac{1}{N\beta}\int p(J_{ij})dJ_{ij} \ln Z_N(\beta, \{J_{ij}\}) \quad \dots (28)$$

The free energy per spin of a coupling realization $\{J_{ij}\}$ of a N spin system is,

$$\Xi_N = -\frac{1}{N\beta}\ln Z_N(\beta, \{J_{ij}\})$$

The self-averaging of Ξ_N is

$$F = \lim_{N\to\infty} \Xi_N$$

For a unidimensional system with first neighbor interactions and uniform field h, we can write the partition function as the trace over 2×2 random transfer matrices product. The Hamiltonian is now $H = -\sum_i (J_i\sigma_i\sigma_{i+1} + h\sigma_i)$, such that,

$$Z_N = Tr\prod_{i=1}^{N} M_i, M_i = \begin{pmatrix} e^{\beta J_i+\beta h} & e^{-\beta J_i+\beta h} \\ e^{-\beta J_i-\beta h} & e^{\beta J_i-\beta h} \end{pmatrix} \quad \ldots(29)$$

The moments of the partition function can be estimated as an integral over the spectrum of the possible free energies $[\Xi_{\min}, \Xi_{\max}]$,

$$\langle Z_N(\beta)^q \rangle \propto \int \prod(\Xi)\, d\Xi\ e^{(-\beta \Xi q N)} \quad \ldots(30)$$

The Kolmogorov entropy is related to the sum of the positive Lyapunov exponents which measure the divergence rate along the expanding directions, in accordance with the Theory of Physical Abstraction. For an ergodic measure with a compact support (as proved by Pesin) is,

$$K_1 \le \sum_{i=1}^{P} \lambda_i\,;$$

where P is the number of exponents, $\lambda_i > 0$. In Hamiltonian systems,

$$K_1 = \sum_{i=1}^{P} \lambda_1 = \left.\frac{dL^{(P)}}{dq}\right|_{q=0}$$

A record of measures of a signal $x(t)$ at uniform spacing τ is

$$x_i = x(i\tau); i = 1,2,\ldots, M \gg 1 \quad \ldots(31)$$

Clustering:

Since the number of eddies at scale l with singularity h is proportional to $l^{-d(h)}$, the number of grid points that have to be considered for resolving the set $S(h)$ is

$$N_h(R_e) \sim (L/\eta(h))^{d(h)} \propto R_e^{d(h)/(1+h)};$$

where $R_e = \frac{(\varepsilon L^4)^{1/3}}{\nu}$ and η is the dissipative Kolmogorov length.

Integrating over h, the total number of degrees of freedom is,

$$N(R_e) = \int d\rho(h) N_h(R_e) \propto R_e^{\delta};$$

where $\delta = \max_h[d(h)/1 + h]$.

The estimate $l_{min} = \eta(h_{min})$ assures that all the sets $S(h)$ are taken into account. The number of equations which allows us to get such a fully accurate description is thus;

$$N_T^* \sim \left(\frac{L}{l_{min}}\right)^3 \propto R_e^{3/(1+h_{min})};$$

which may be obtained by considering flows in the required number of directions or dimensions.

Let us define an effective mass dimension $\widetilde{D}$ of the point on which the energy dissipation is concentrated by,

$$\tilde{h} = \frac{(\widetilde{D} - 2)}{3}.$$

$D_1 \simeq 2.87$, which corresponds to select a $\tilde{h} = d\zeta_P/dP|_{P=3} \simeq 0.29$. $\tilde{l}$ is the smallest scale on which average active eddies are still present.

The minimum separation Δp between disturbances, with difference $\Delta\lambda$, is such that

$$\Delta p = Q\Delta\lambda \quad \dots (32);$$

$Q = 0.00642986 \times 10^5$ from experimental results.

Thus any number of points n inside a given representation will form a cluster point if they sufficiently close, as described by equation (32). The total information inside such a cluster is,

$$I_\varepsilon = \frac{1}{n}\sum_{i=1}^{n} I_{\varepsilon_i} \quad ..(33);$$

where I_{ε_i} is the energy dissipation information inside the constituent points. As the constituent points must be sufficiently close in order to form a cluster point, they may be considered to be a continuous energy dissipation information function, such that,

$$I_\varepsilon = \frac{1}{n}\int_1^n I_\varepsilon \, d\varepsilon$$

As I_ε is a function of ε, we can write,

$$\int_1^n I_\varepsilon \, d\varepsilon = \lim_{h\to 0} h \sum_{r=1}^{n} f(1 + rh) \quad \dots (34);$$

where $nh = n - 1$, such that, $h = \frac{n-1}{n}$.

As $n \to \infty, h \to 0$.

A corollary to the Laws of Physical Transactions suggest that the individual constituent points of the cluster point will tend to be in the lowest effective dissipation energy state $\tilde{\varepsilon}$. Thus, the total energy dissipation information that the cluster point will tend to reach is,

$$\widetilde{I_\varepsilon} = nI_{\tilde{\varepsilon}}$$

The total loss in energy dissipation energy information by the cluster point is therefore,

$$\Delta I_\varepsilon = I_\varepsilon - \widetilde{I_\varepsilon} = \frac{1}{n}\sum_{i=1}^{n} I_{\varepsilon_i} - I_{\tilde{\varepsilon}}$$

Let the individual energy dissipation information states of the constituent points be,

$$I_{\varepsilon_i} = \alpha_i I_{\tilde{\varepsilon}},$$

such that the loss in dissipation energy information is,

$$\Delta I_{\varepsilon} = I_{\varepsilon} - \widetilde{I_{\varepsilon}} = I_{\tilde{\varepsilon}}\left[\left(\frac{1}{n}\sum_{i=1}^{n}\alpha_i\right) - 1\right] \qquad \dots (35)$$

Taking a continuous distribution,

$$\Delta I_{\varepsilon} = I_{\tilde{\varepsilon}}\left[\frac{1}{n}\left(\int_{1}^{n}\alpha\, d\alpha\right) - 1\right]$$

$$i.e., \Delta I_{\varepsilon} = I_{\tilde{\varepsilon}}\left[n\left(\frac{n^2-1}{2}\right) - 1\right] \qquad \dots (36).$$

A close enough representation of the constituents of the universe therefore will be affected by this loss in dissipation energy information ΔI_{ε} inside its various clusters. This in turn will give rise to seemingly anomalous behaviour inside the representation. The clusters will seem to move away from each other with greater velocities than anticipated values. On the other hand, the clusters themselves will seem to be bound with greater strengths than is anticipated. The existence of dark energy and dark matter that we feel may be linked to the loss ΔI_{ε}.

Conclusion:

Choosing a suitable scaling ratio we may represent the microscopic and the macroscopic worlds on the same scale, enabling us to study and compare their various hitherto hidden properties. A large enough packing density ensures formation of cluster points inside the representations. These cluster points will tend to be in their lowest energy dissipation states. The whole universe being considered as a cluster when its constituents are close enough, as in the moment of the Big Bang, it tends to be in its lowest energy state (theoretically a zero energy state). ΔI_{ε}, i.e., the difference in dissipation energy information which tends to infinity as the number of constituent points inside it tends to infinity, however establishes itself as Big Bang takes place. Yet, as the universe can expand further as its constituents move away from each other, with respect to a further expanded state it at any given present representation its clustered. Thus a hidden amount of energy dissipation information is present at any moment we look at the universe. This hidden energy dissipation information will make the clusters to move away from each other and the clusters themselves to be bound within themselves with greater hidden strengths than is anticipated.