Chapter 7

Gravitational fields of the topological singularities

Thanks to *Newton's second partial equation*, one can obtain *the external fields of expansion perturbations*, i.e. *the external fields of gravitation* associated with a localized topological singularity. We find that these external gravitational fields can have three components, namely a generally dominant component associated with *the elastic energy of distortion, that is to say with the mass of the singularity*, and two components generally much weaker due respectively to *the charge of curvature* of the singularity and to *the charge of rotation* of the singularity.

In the process, we also come to show that the collapse of clusters of singularities of the vacancy type or of the interstitial type leads to fairly singular macroscopic topological singularities within the lattice: a hole in the lattice, a kind of *macroscopic vacancy*, in the case of the collapse of singularities of vacancy nature, or a piece of additional lattice, a kind of *macroscopic interstitial* within the lattice in the case of the collapse of singularities of interstitial nature. The description of the “gravitational” fields of these two types of macroscopic complementary singularities shows that the macroscopic vacancy singularity can behave like a real black hole, whereas the macroscopic interstitial singularity does not have this property. Subsequently, these macroscopic singularities will prove to be ideal candidates to explain the *"black holes"* of the Universe in the case of the macroscopic vacancies and the *"neutron stars"* of the Universe in the case of the macroscopic interstitials.

By applying the calculations of the external gravitational field of the topological singularities to the microscopic singularities in the form of twist disclination loops, prismatic edge dislocation loops or slip mixed dislocation loops, we deduce the set of gravitational properties of these loops. Several extremely interesting consequences will be deduced therefrom, in particular the existence, in the case of the prismatic edge dislocation loop, of an equivalent mass of gravitation different from the mass of inertia, which may even prove to be negative in the case of loops of an interstitial nature, a result which will have very significant consequences thereafter.

Expansion perturbations by a singularity with a given distorsion energy

Let us consider a singularity localized at rest, of volume $\boldsymbol{V}_{cluster}$, consisting of a loop or a cluster of numerous loops of dislocations and / or of disclinations, and let us suppose known the densities of distortion energy $F_{dist}^{cluster}(\vec{r})$ and potential energy $F_{pot}^{cluster}(\vec{r})$ within this singularity. The equilibrium of the expansion perturbation field $\tau_{int}^{(E)}(\vec{r})$ within this singularity is given by the solution of the second degree equation deduced from Newton's second partial equation in figure 5.1. We cannot do here an exact calculation of $\tau_{int}^{(E)}(\vec{r})$ since it would require dealing with a concrete case of singularity to know exactly the distributions of the densities of distortion energy $F_{dist}^{cluster}(\vec{r})$ and potential energy $F_{pot}^{cluster}(\vec{r})$ within this singularity. On the other hand, one can deal with this problem in an approximate way by introducing average values of the various fields concerned. We can start by introducing *the global energy at rest* $E_{dist}^{cluster}+V_{pot}^{cluster}$ of the

topological singularity by summing the distributions of the densities of distortion energy $F_{dist}^{cluster}(\vec{r})$ and potential energy $F_{pot}^{cluster}(\vec{r})$ within this singularity, so that we can define an average value $\overline{\tau}_{int}^{(E)}$ of the internal field of expansion perturbations. It is clear that this mean field is purely virtual, in other words that it does not really exist, but that it represents in fact a form of mean value of all the accidents of the perturbation field within the singularity, accidents which must be extremely marked especially if we have to deal with a cluster of very many topological singularities. The exact calculation of the mean perturbation field $\overline{\tau}_{int}^{(E)}$ is quite complicated, since it is subject to various assumptions regarding its conditions of existence. We will not enter here into a detailed discussion of these conditions of existence. The detailed calculations show that there are static or dynamic solutions for the expansion disturbance fields associated with the topological singularity, as shown in Figure 7.0.

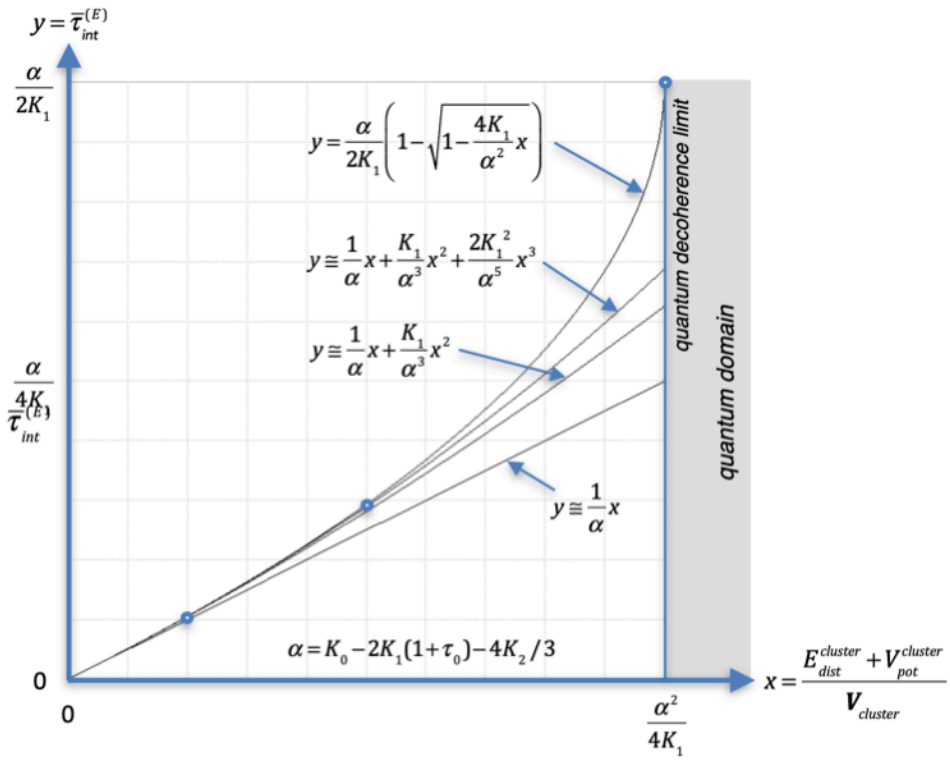


Figure 7.0 *- the function $\overline{\tau}_{int}^{(E)}$ in the static solution domain for energy density values of the singularity below the critical value $\alpha^2/4K_1$, and domain of the quantum behaviour for values above the critical value.*

The expansion field perturbations associated with a localized topological singularity are in fact an expression of the existence of *a static external "gravitational field" at a long distance from this singularity*, as long as it has an energy density or rotational load density below a certain critical value $\alpha^2/4K_1$.

In the case where this energy density becomes greater than this critical value $\alpha^2/4K_1$, the expansion field associated with this localized topological singularity becomes *a dynamic expansion perturbation, which will bring to light the quantum behaviors of this singularity,* which we will discuss in chapter 11. The critical value of the energy density then becomes an extremely important quantity since it corresponds in fact to *a quantitative value which defines*

the famous quantum decoherence limit, i.e. the limit of passage between a classical behaviour and a quantum behaviour of the topological singularity.

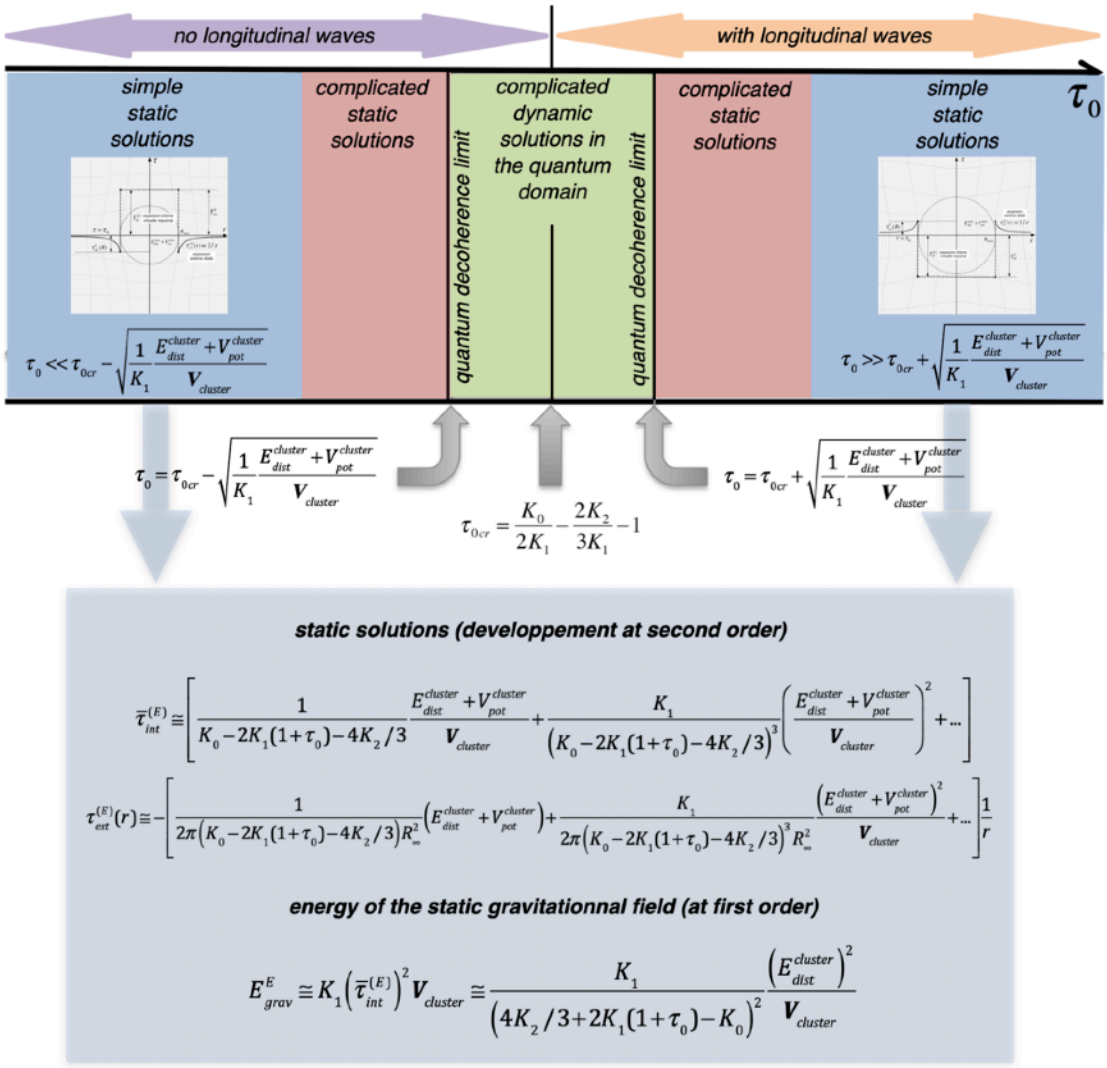


Figure 7.1 *- the domains of solutions of expansion perturbations of a singularity with a given energy* $E_{dist}^{cluster}+V_{pot}^{cluster}$ *as a function of the backgroud expansion*

The background expansion domains of the lattice in which these solutions occur are given by conditions which are shown in figure 7.1. In a small domain centered on $\tau_0 = \tau_{0cr}$, there can only exist a dynamic solution of Newton's equation for the expansion perturbation, dynamic solution which allows to pass from the static solution of the domain $\tau_0 < \tau_{0cr}$ to the static solution in the domain $\tau_0 > \tau_{0cr}$ and vice-versa.

In the domains outside this small domain of dynamic solutions, there are static solutions for the perturbations of the volume expansion fields caused by the singularity. But here too, the calculations and the results obtained are quite complicated. But there are areas of lattice background expansion τ_0 that present simple static solutions. These domains are reported in figure 7.1, with the expressions of the internal and external perturbation fields which correspond to them. It is remarkable to note then that the average value $\overline{\tau}_{int}^{(E)}$ of the internal field of

perturbation of the expansion depends on *the global energy of rest* $E_{dist}^{cluster}+V_{pot}^{cluster}$ of the singularity and its volume $V_{cluster}$, while the field of external perturbations decreases in $1/r$ with the distance from the cluster, that it only depends on the total energy $E_{dist}^{cluster}+V_{pot}^{cluster}$ of it, and especially that *it does not depend on the volume or the radius of it.*

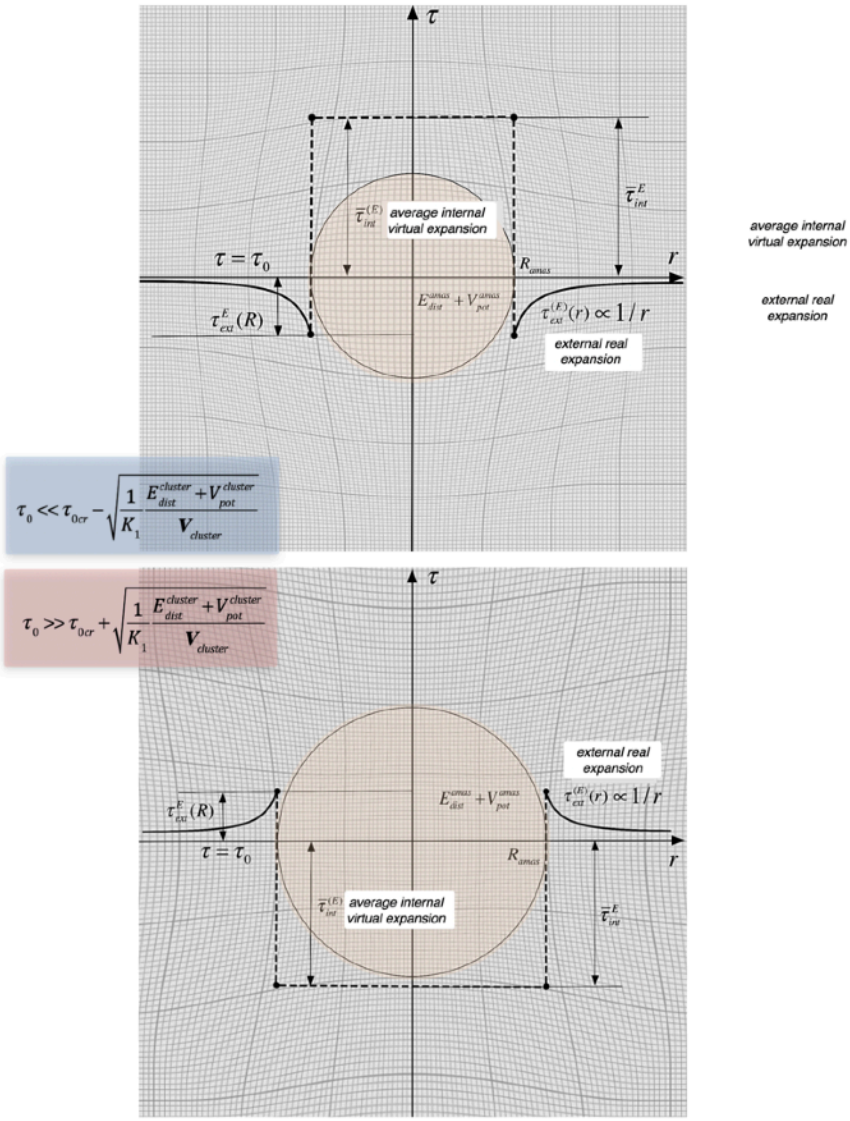


Figure 7.2 *- the expansion field inside and outside a singularity with a given energy* $E_{dist}^{cluster}+V_{pot}^{cluster}$

These behaviors of the expansion perturbation field are reported in figure 7.2 as a function of the distance r from the center of the singularity, in both cases $\tau_0 << \tau_{0cr}$ and $\tau_0 >> \tau_{0cr}$.

The total energy of the expansion field can also be calculated, and we see that it is positive, that it is proportional to the product of the volume of the singularity by the square of the internal field of expansion perturbations of the singularity, and that it depends only on the modulus of elasticity K_1 which, as we recall, must be much smaller than the modules K_0 and K_3 in the perfect cosmological lattice.

It is interesting to compare the gravitational energy of the singularity with its elastic energy of distortion, by calculating the approximate value of the ratio of the two. The expression $\boldsymbol{E}_{grav}^{E} / \boldsymbol{E}_{dist}^{cluster} \cong (K_1 / K_0^2)(\boldsymbol{E}_{dist}^{cluster} / V_{cluster})$ comes for $\tau_0 << \tau_{0cr}$. Now, according to conjecture 6, the ratio K_1 / K_0^2 is extremely small, so that the gravitational energy of the singularity is certainly extremely smaller than the elastic energy of distortion of the singularity in the domain $\tau_0 << \tau_{0cr}$. We will come back to this point later.

With the conjectures 6 of the *perfect cosmological lattice*, if we are in the *expansion domain* $\tau_0 < \tau_{0cr}$ in which there are no longitudinal waves, the mean virtual internal field $\overline{\tau}_{int}^{(E)}$ of the singularity is positive. As for the real external field $\tau_{ext}^{(E)}(\vec{r})$ of the singularity, it is negative and therefore perfectly satisfies conjecture 3 of figure 3.4 deduced from the curvature of the wave rays in the vicinity of the singularity.

On the other hand, if one is in the *domain of expansion* $\tau_0 > \tau_{0cr}$ where there are longitudinal waves, the fields reported in figure 7.2 are reversed compared to the fields in the domain of expansion $\tau_0 < \tau_{0cr}$: the average virtual internal field $\overline{\tau}_{int}^{(E)}$ of the singularity becomes negative, and the real external field $\tau_{ext}^{(E)}(\vec{r})$ becomes positive, so that it no longer satisfies conjecture 2 deduced from the curvature of the wave rays in the vicinity of the singularity.

Expansion perturbations by a singularity with a given curvature charge

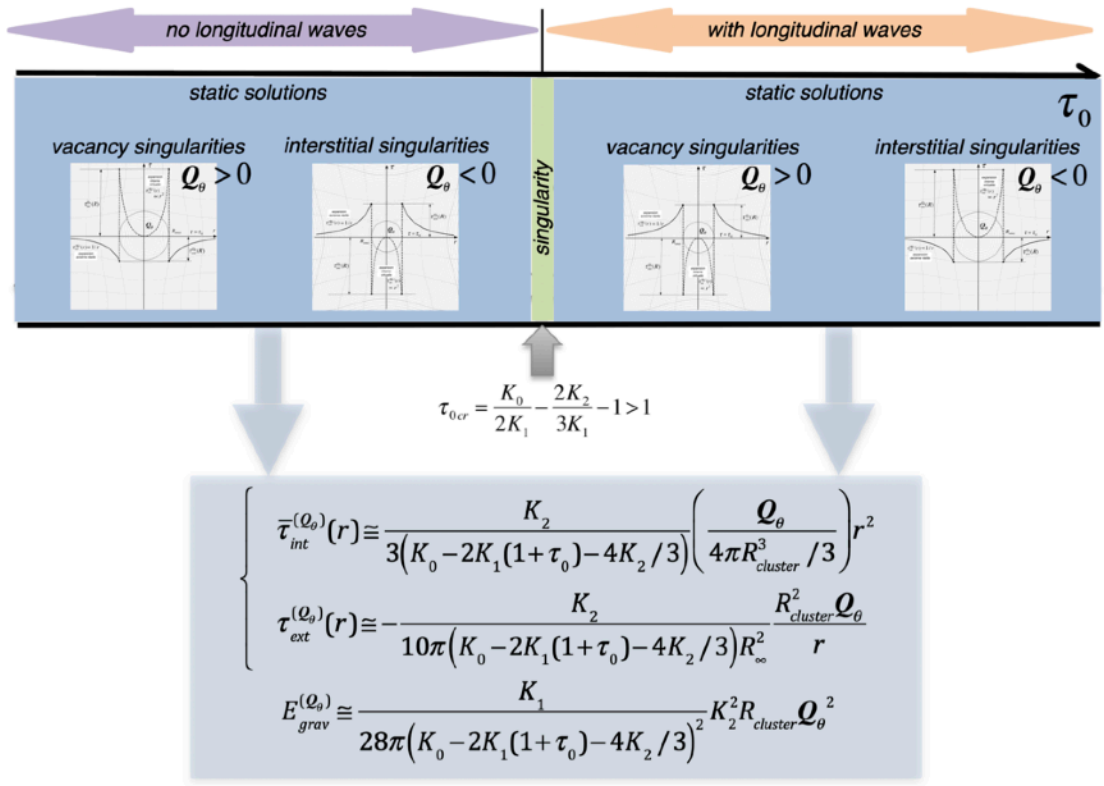


Figure 7.3 *- the domains of solutions of expansion perturbations of a singularity with a given curvature charge* Q_θ *as a function of the backgroud expansion*

A localized singularity of radius $R_{cluster}$, apart from having a given rest energy, can also have an overall charge of curvature Q_θ. Indeed, such a singularity can consist of a cluster of discrete topological singularities of the lattice, such as prismatic dislocation loops which each have a n elementary charge of curvature $q_{\theta BC(i)}$. If $Q_\theta > 0$, we will speak of *a singularity of a vacancy nature* because there are missing laattice sites within the cluster, and if $Q_\theta < 0$, we will speak of *a singularity of interstitial nature*, because there is then excess of lattice sites within the cluster. A curvature singularity is responsible for a non-zero and divergent flexion field in its vicinity as we have shown in figure 2.33.

To find the perturbations of the expansion field due to this singularity, we start from the static equilibrium equation obtained by the divergence of Newton's first partial equation in figure 5.1. The equilibrium equation of the expansion field of such a singularity then requires knowing the density $\theta^{cluster}(\vec{r})$ of curvature charges within the singularity. Generally not knowing this density which is specific to a given singularity, we use to simplify the approximation of a homogeneous

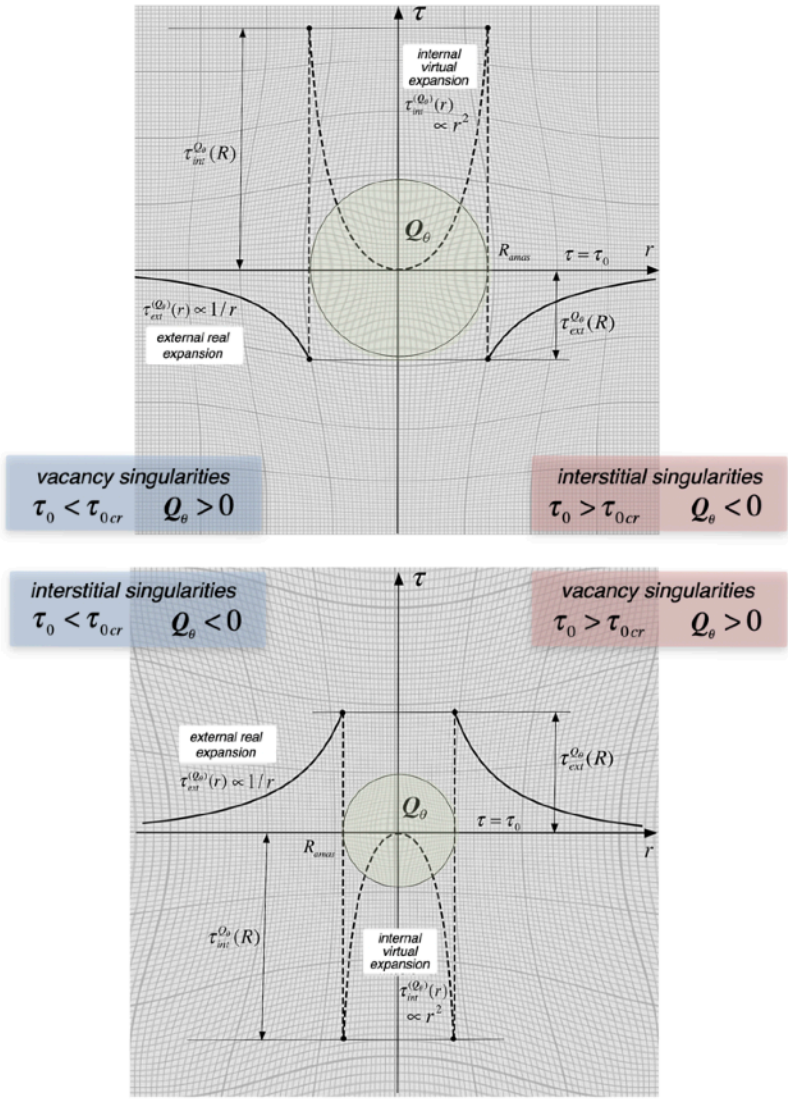


Figure 7.4 *- the expansion field inside and outside a singularity with a given curvature charge* Q_θ

average curvature charge density within the singularity, such as $\overline{\theta}^{cluster} = \boldsymbol{Q}_\theta / \left(4\pi R_{cluster}^3 / 3\right)$. We then obtain a simplified equation of equilibrium for the virtual internal field whose solution $\tau_{int}^{(\boldsymbol{Q}_\theta)}(r)$ is of spherical symmetry. As before, the exact calculation of the perturbation fields is quite complicated, because it is subject to various assumptions concerning its conditions of existence, in particular the condition that the total internal and external perturbation field is of zero mean value, so that the value of the mean field within the lattice is equal to the background expansion τ_0 of the lattice. Under these conditions, we obtain the internal and external fields reported in figure 7.3, as well as the energy associated with this expansion perturbation field.

These expressions imply that, if these fields exist, that is to say if $K_2 \neq 0$, these present an infinite singularity when the background expansion τ_0 reaches the critical value τ_{0cr} which cancels the denominator, and we can show that a singularity of curvature necessarily goes through a black hole stage, for $\tau_0 \leq \tau_{0cr}$ in the case of singularities of a vacancy nature and for $\tau_0 \geq \tau_{0cr}$ in the case of singularities of interstitial nature.

The signs of the term $K_0 - 2K_1(1+\tau_0) - 4K_2/3$ and of the module K_2 play an important role here. Using the conjectures 6 of the *perfect cosmological lattice*, we have that:

- if $K_2 / \left[K_0 - 2K_1(1+\tau_0) - 4K_2/3\right] > 0$, that is to say if $\tau_0 < \tau_{0cr}$, *only the singularities of a vacancy nature*, with $\boldsymbol{Q}_\theta > 0$, satisfy the conjecture 3 of figure 3.4 concerning the curvature by attraction of the wave rays in the vicinity of the singularity, while *the singularities of interstitial nature*, with $\boldsymbol{Q}_\theta < 0$ repel the wave rays,
- if $K_2 / \left[K_0 - 2K_1(1+\tau_0) - 4K_2/3\right] < 0$, that is to say if $\tau_0 > \tau_{0cr}$, only the singularities of interstitial nature, with $\boldsymbol{Q}_\theta < 0$, satisfy conjecture 3 concerning the curvature by attraction of the wave rays in the vicinity of the singularity, while the singularities of vacancy nature, with $\boldsymbol{Q}_\theta > 0$, repel the wave rays.

These behaviors of the expansion perturbation field are reported in figure 7.4 as a function of the distance r from the center of the vacancy or interstitial singularity, in the four possible cases, namely for $\tau_0 < \tau_{0cr}$ and $\tau_0 > \tau_{0cr}$, and for $\boldsymbol{Q}_\theta < 0$ and $\boldsymbol{Q}_\theta > 0$.

Expansion perturbations by a singularity with a given rotation charge

Now imagine the existence within the perfect lattice of a localized singularity of volume $\boldsymbol{V}_{cluster}$ and global rotation charge $\boldsymbol{Q}_\lambda$, composed of a cluster of elementary rotation charges $\boldsymbol{q}_{\lambda BV(i)}$ or containing a density $\lambda(\vec{r})$ of rotation charges.

Such a singularity will therefore have an external field of expansion perturbations linked to the elastic energy of the singularity due to the field of internal rotation at the singularity, and which is deduced directly by the second relation of figure 7.1. This perturbation field therefore does not depend on volume $\boldsymbol{V}_{cluster}$ or radius $R_{cluster}$ of the singularity with rotation charge.

But to this external field of expansion perturbation due to the internal energy of rotation of the singularity must also be added an external field of expansion perturbations which is generated by the external field of rotation of the singularity, and which appears as a result of the distortion energy $E_{dist(rot\,ext)}^{cluster}$ associated with the external rotation field $\vec{\omega}_{ext}^{él}$. The external rotation field of the singularity is given in figure 5.4, and it does not depend either on the volume $\boldsymbol{V}_{amas}$ nor on the radius $R_{cluster}$ of the rotation singularity. This external perturbation $\tau_{ext}^{(E_{dist(rot\,ext)})}(r)$ of the volume expansion field is deduced using the static version of Newton's second partial equation of figure 5.1, in which the energy density is due to the external elastic energy of the rotation field

of the singularity, and in which τ^{ext} and τ^{ch} are zero. The solution of the second degree equation thus obtained presents several solutions depending on the volume expansion domain τ_0 considered, as shown in figure 7.5

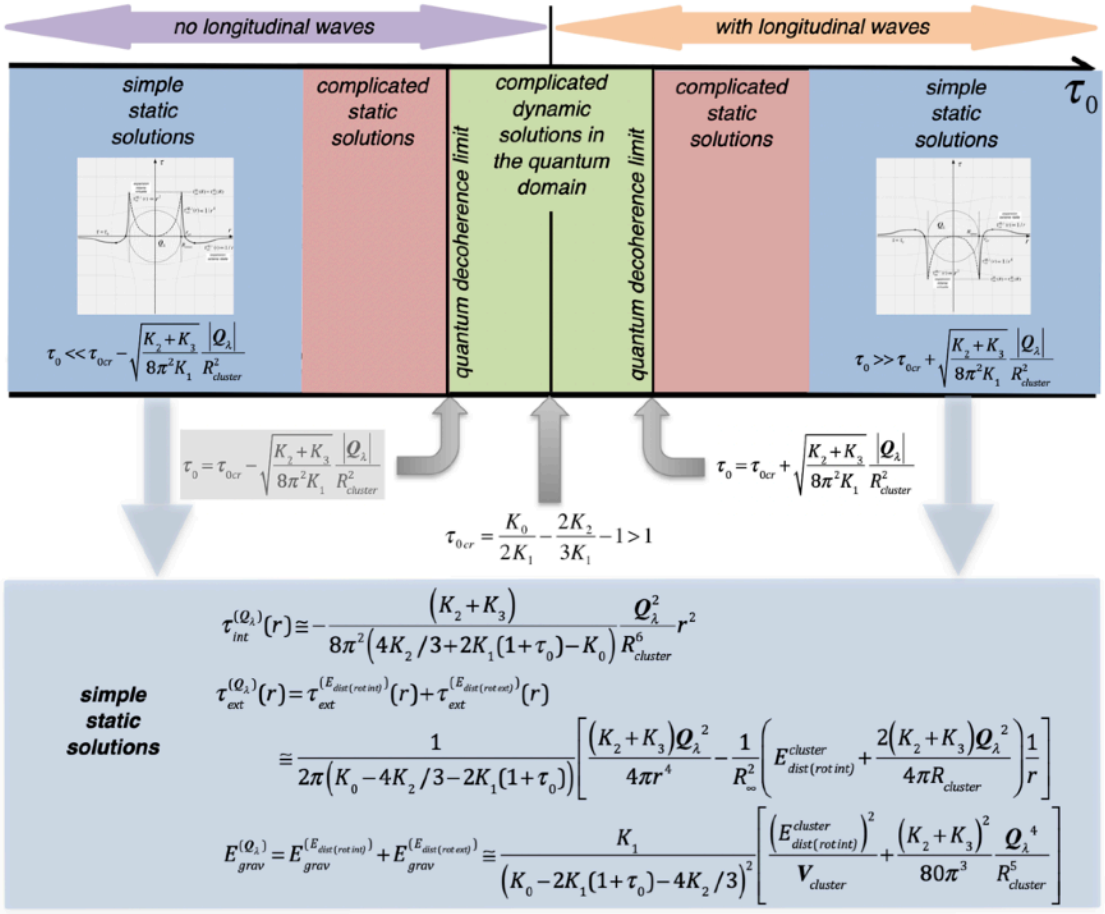


Figure 7.5 *- the domains of solutions of expansion perturbations of a singularity with a given rotation charge* Q_λ *as a function of the backgroud expansion*

In a domain of volume expansion τ_0 centered on $\tau_0 = \tau_{0cr}$, there can only exist *a dynamic solution of Newton's equation for the perturbation of external expansion associated with the singularity of rotation,* which makes it possible to pass from the static solution of the domain $\tau_0 < \tau_{0cr}$ to the static solution in the domain $\tau_0 > \tau_{0cr}$ and vice versa. The transition limit from the dynamic solution to the static solution again corresponds to *the quantum decoherence limit.*

In the vicinity of this domain of dynamic solutions, there appear two domains of fairly complicated static solutions which will not be discussed here. Finally, it is for the extreme values of the background field τ_0 that fairly simple static solutions appear for the external field of perturbations due to the singularity of rotation. Once again assuming that the mean value of the expansion taken outside and inside the singularity is equal to τ_0, we obtain the final expression of the external perturbation field $\tau^{(Q_\lambda)}_{ext}(r)$ reported in figure 7.5.

It can therefore be seen that the total external field of expansion perturbations of a singularity of rotation has a long-range component in $1/r$ associated with the internal and external elastic

energy of rotation of the singularity and a short-range component in $1/r^4$ associated directly with the energy of the external rotation field of the singularity. We superimposed very schematically the superposition of these two fields in figure 7.6, respectively in the cases where $\tau_0 < \tau_{0cr}$ and $\tau_0 > \tau_{0cr}$.

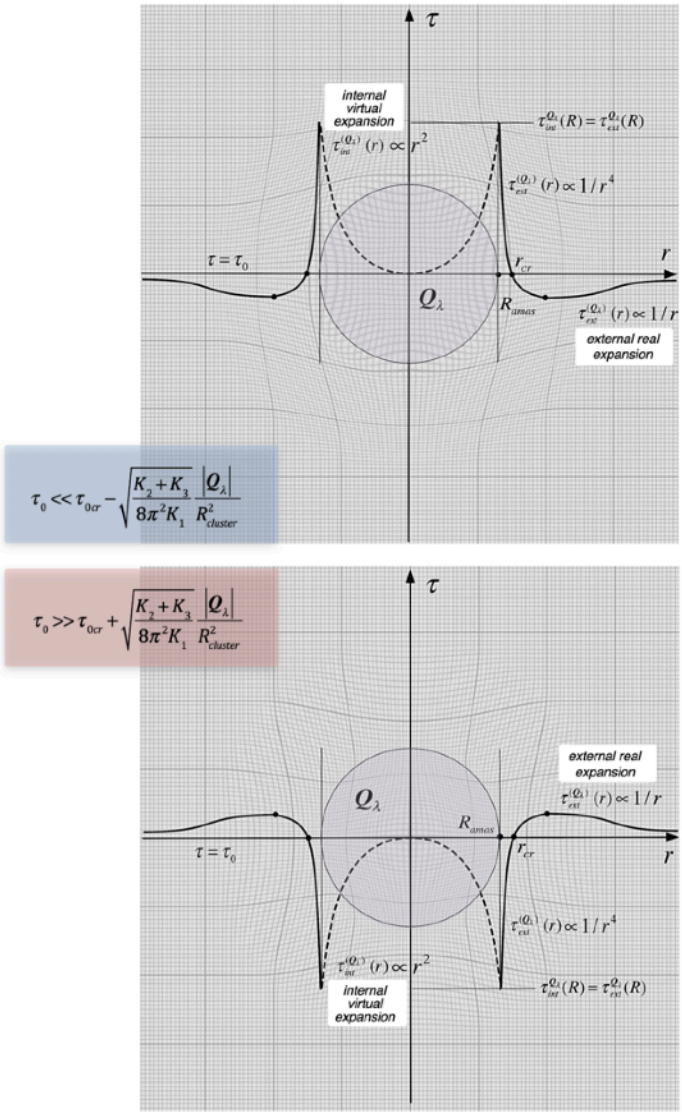


***Figure 7.6** - the expansion field inside and outside a singularity with a given rotation charge Q_λ*

In both cases (spherical singularity of rotation or twist disclination loop), it can be shown that there is a critical distance r_{cr} greater than the radius of the singularity $(r_{cr} > R_{cluster})$ for which the external expansion field is reversed. We also note that the perturbations of the expansion field at the interface of the singularity are positive in the expansion domain $\tau_0 < \tau_{0cr}$ and negative in the expansion domain $\tau_0 > \tau_{0cr}$.

On the other hand, as the energies of distortion (for example $E_{dist}^{(Q_\lambda)}$ in the case of a spherical load of rotation or $E_{dist\,tore}^{BV}$ in the case of a twist disclination loop) are always positive, there is no

asymmetry in the fields of expansion between charges and anti-charges of rotation. On the other hand, there appears to be an inversion of these fields when moving from the domain $\tau_0 < \tau_{0cr}$ to the domain $\tau_0 > \tau_{0cr}$.

In figure 7.5, one also reported the field of the internal perturbation of expansion, in the case of a spherical charge of rotation for which this field can be calculated in an exact way. We also reported the total energy of the gravitational perturbation field, also calculated in the case of a spherical charge of rotation.

About the gravitational fields of macroscopic singularities

Concerning the fields of "gravitational" nature of the macroscopic singularities which we have just described, it is very encouraging to note that there is a first field of expansion associated directly with *the energy* $E_{dist}^{cluster} + V_{pot}^{cluster}$ *of the cluster of topological singularities*, in a similar way to the gravitational field of Einstein's General Relativity which is also an emanation of *the energy-momentum tensor of matter.*

But there still appear in our approach two other "gravitational" fields of expansion which are associated respectively with *the global charge* $\boldsymbol{Q}_\theta$ *of curvature* and *the global charge* $\boldsymbol{Q}_\lambda$ *of rotation* of the cluster of topological singularities. These fields have in fact no equivalent in Einstein's General Relativity.

The existence of the second "gravitational" expansion field, due to *the curvature charge* $\boldsymbol{Q}_\theta$, is subject to the condition that the shear modulus K_2 of the perfect lattice is not zero. There is therefore still the possibility of discussing the existence or not of this field according to the value which must be attributed to the module K_2 in our analogy with the real world, knowing that this module must in any case be very small vis-à-vis of the module K_3, as already specified with conjecture 4.

The third "gravitational" expansion field is associated with *the rotation charge* $\boldsymbol{Q}_\lambda$ of the cluster of topological singularities considered. Within the framework of our analogy, this third expansion field must necessarily exist if the cluster has a non-zero charge $\boldsymbol{Q}_\lambda$ since the module K_3 must exist to satisfy the analogy with Maxwell's equations. But this field has no direct analogy in Einstein's General Relativity and Particle Physics theories.

Note again that the three preceding "gravitational" fields have non-zero gravitational energies. As these depend on the coefficient $K_1 / K_0^2 <<< 1$, which must be very small in the perfect cosmological lattice, the gravitational energy of the singularities is certainly negligible compared to the elastic energy of distortion of the singularities.

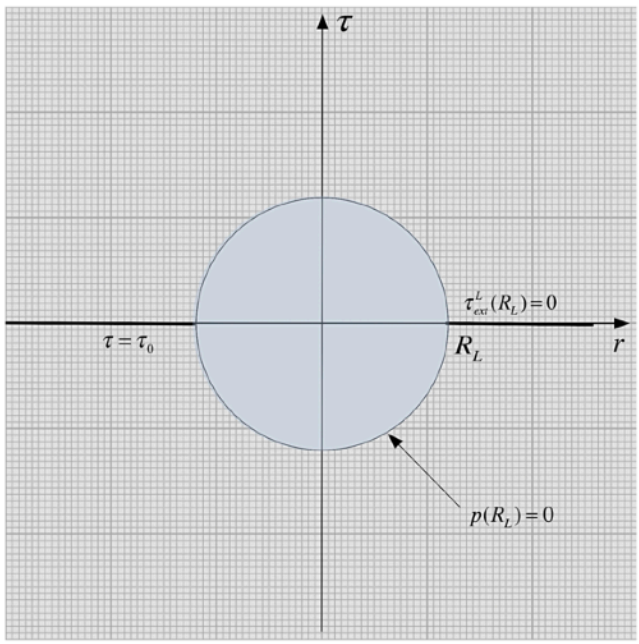
The fact that there appear two "gravitational" fields of volume expansion which apparently have no analogues in the theories of General Relativity of Einstein and in Particle Physics is very interesting to test our approach. We will return in the following chapters to the roles that the three fields of volume expansion could play respectively associated with the energies $E_{dist}^{cluster} + V_{pot}^{cluster}$ the charge of curvature $\boldsymbol{Q}_\theta$ and the charge of rotation $\boldsymbol{Q}_\lambda$ of a localized singularity or of a localized cluster of topological singularities.

If the background expansion field τ_0 of the lattice increases or decreases so as to pass through the critical value τ_{0cr}, there appears a domain where there can no longer exist a static solution to Newton's equation, which means that *it must necessarily appear a dynamic solution which makes it possible to transform a static singularity of the domain* $\tau_0 < \tau_{0cr}$ *into a static*

singularity of the domain $\tau_0 > \tau_{0cr}$ *, and vice versa*. The non-existence of a static solution to Newton's equation when the energy density and / or the rotational charge of the singularity becomes too large is mathematically a fact quite similar to what we already encountered in chapter 2 in the case of *Frank-Read sources of dislocations*, where there were no longer any static solutions to the deformation of an anchor dislocation string when the stress exceeded a certain critical limit value, and we will see below that the appearance of pure dynamic solutions for the expansion perturbation field have in fact a very close link with quantum physics.

Macroscopic vacancy located in the lattice, real gravitational black hole

Imagine that a cluster of vacancy type singularities, i.e. singularities carrying positive curvature charges, such as prismatic dislocation loops of the lacunar type for example, collapses on itself (under the effect for example of "gravitational attractive forces" which we will



$$\begin{cases} \tau_{ext}^{(L)}(r) \cong -\left(1+\tau_0+\tau^{external}(R_L)\right)\dfrac{R_L}{r} \\ R_L = \sqrt[3]{\dfrac{3N_L}{4\pi n_0 \mathrm{e}}} \end{cases}$$

$$E_{grav}^{(L)} \cong 2\pi\left(K_0-2K_1\tau_0\right)\left(1+\tau_0+\tau^{external}(R_L)\right)R_\infty^2 R_L$$

$$\cong \sqrt[3]{\frac{6\pi^2 N_L}{n_0 \mathrm{e}}}\left(K_0-2K_1\tau_0\right)\left(1+\tau_0+\tau^{external}(R_L)\right)$$

Figure 7.7a *- the expansion field* $\tau_{ext}^{(L)}(r)$ *of a macroscopic vacancy of about 10'000 lattice sites for* $\tau_0 = -\left(1+\tau^{ext}(R_L)\right)$*, and the expressions of its radius and its gravitational energy*

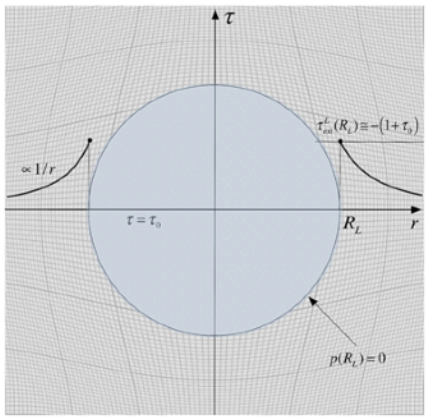


Figure 7.7b *- the expansion field $\tau^{(L)}_{ext}(r)$ of a macroscopic vacancy of about 10'000 lattice sites for $\tau_0 < -\left(1+\tau^{externe}(R_L)\right)$, represented on the same scale as in figure 7.7a*

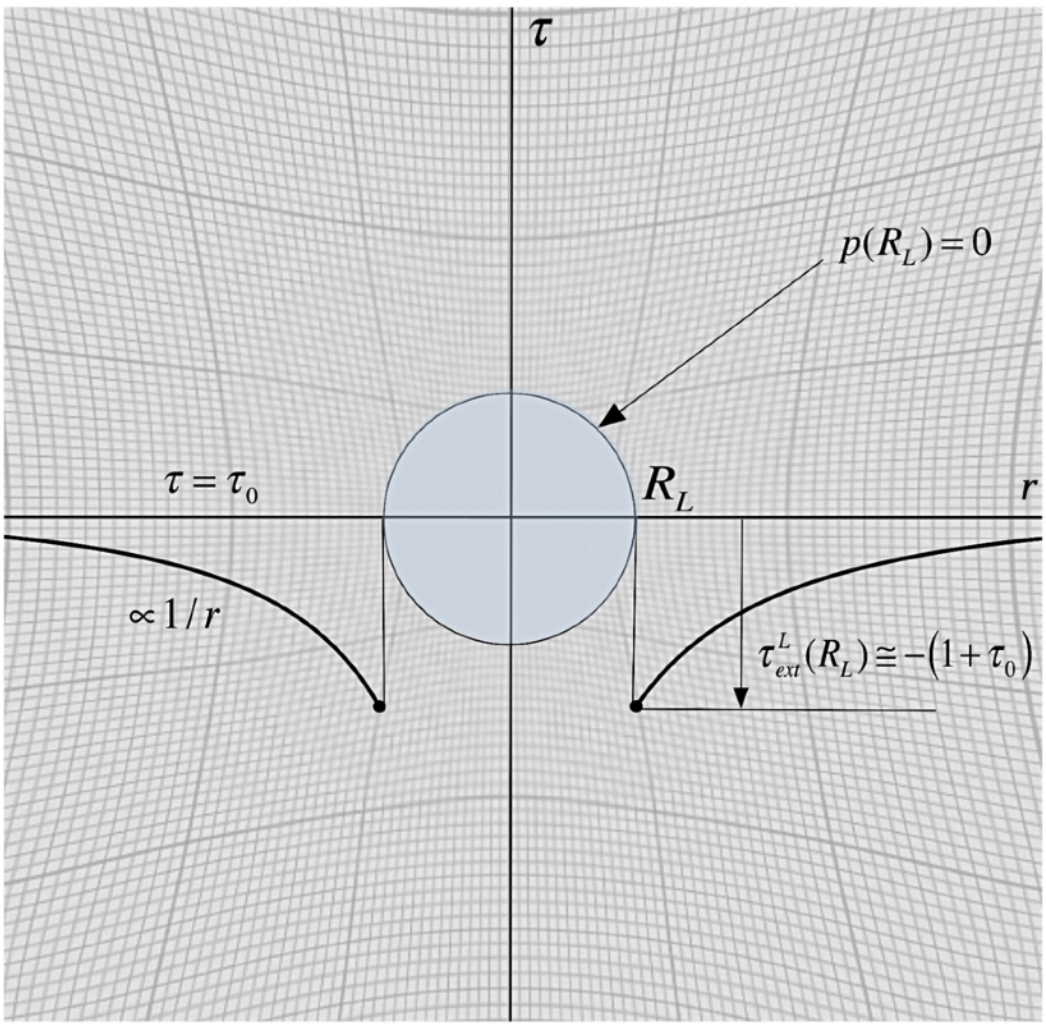


Figure 7.7c *- the expansion field $\tau^{(L)}_{ext}(r)$ of a macroscopic vacancy of about 10'000 lattice sites for $\tau_0 > -\left(1+\tau^{externe}(R_L)\right)$, represented on the same scale as in figure 7.7a*

describe later). If the initial cluster is neutral vis-à-vis the rotation charges, the individual singularities of the cluster combine, losing their own identity (dislocation or disclination loops) to form a single macroscopic hole within the lattice, as a kind of macroscopic vacancy formed by individual vacancies, as shown in figures 7.7a to 7.7c. This means that there is a local lack of N_L lattice sites. The radius of this macroscopic vacancy, assumed to be spherical, is then worth $R_{L0} = \sqrt[3]{3N_L / (4\pi n_0)}$, in the imaginary case where the lattice presented a homogeneous expansion of $\tau = 0$.

Within the real lattice, the presence of this macroscopic vacancy will generate a spherical volume expansion field $\tau_{ext}^{(L)}(r)$ that must be defined. On the surface of the singularity, the expansion field is equal to the sum of the field $\tau_{ext}^{(L)}(R_L)$ generated by the singularity, the background expansion field τ_0 of the lattice and an external expansion field $\tau^{external}(R_L)$ due to the other singularities located in the vicinity of the macroscopic vacancy. On the other hand, on the surface of the singularity, the total field must be arranged so that the pressure at the interface of the hole is zero. To obtain *the field of expansion* $\tau_{ext}^{(L)}(r)$ of the macroscopic vacancy as well as *its real radius* R_L, the condition of zero pressure at the interface can be injected into the equation of state of the pressure in figure 3.1. We obtain a second degree equation, which, under the assumption $K_0 >> K_1$ of conjecture 6, has only one valid solution which is shown in figure 7.7a. We can then show that the mean expansion field in the presence of this macroscopic singularity is indeed equal to the background expansion τ_0 of the lattice.

On the other hand, we note that the real radius R_L of the macroscopic vacancy is a quantity which depends only on the number N_L of missing lattice sites.

In figures 7.7a to 7.7c, the expansion field associated to a macroscopic vacancy of approximately 10,000 substitutional sites has been plotted graphically, for three different values of the background expansion τ_0 of the lattice, and trying to respect at best the scale of the drawing.

The energy of a macroscopic vacancy associated with its field $\tau_{ext}^{(L)}(r)$ of expansion perturbations is calculated quite easily, and we obtain the relation reported in figure 7.7a. It can be seen that the energy of formation of the macroscopic vacancy depends only on the background expansion τ_0 of the lattice and that it is canceled $(\boldsymbol{E}_{grav}^{(L)} \to 0)$ for the values $\tau_0 \to -\left(1+\tau^{external}(R_L)\right)$ and $\tau_0 \to K_0/(2K_1)$. Between these two values, namely in the interval $-\left(1+\tau^{external}(R_L)\right) < \tau_0 < K_0/2K_1$, the energy of formation of the macroscopic vacancy is positive, whereas it becomes negative outside this interval. On the other hand, within the interval $-\left(1+\tau^{externe}(R_L)\right) < \tau_0 < K_0/2K_1$, the energy of formation of two macroscopic vacancies of N_L sites is higher than the energy of formation of a single macroscopic vacancy of $2N_L$ sites since $\boldsymbol{E}_{grav}^{(L)}\Big|_{2N_L} \cong 2\boldsymbol{E}_{grav}^{(L)}\Big|_{N_L} / \sqrt[3]{4} < 2\boldsymbol{E}_{grav}^{(L)}\Big|_{N_L}$, so that two macroscopic vacancies will have an energy advantage to merge when the background expansion field is in this interval.

In the presence of a macroscopic vacancy, the first condition reported in figure 3.5 for the appearance of a black hole implies that the critical radius to make a black hole with a macrosopic vacancy is equal to $r_{cr} = \left(1+\tau_0+\tau^{external}(R_L)\right)R_L/2$. If we add the necessary condition that $r_{cr} > R_L$, there comes the condition $\tau_0 > 1-\tau^{external}(R_L)$ on the background expansion τ_0 of the lattice to form a black hole.

We conclude that a macroscopic vacancy, whatever its size and its energy, *necessarily becomes a black hole as soon as the background volume expansion* τ_0 *exceeds the value of*

$1-\tau^{external}(R_L)$. This conclusion is extremely interesting insofar as we have here the macroscopic stable topological singularity which we can consider as *the true black hole*, analogous to the black hole of general relativity, when $\tau_0 > 1-\tau^{external}(R_L)$, represented in figure 7.7c. And this vacancy topological singularity behaves like a white hole which repels the wave rays when the condition $\tau_0 < -\left(1+\tau^{external}(R_L)\right)$ is satisfied, as in figure 7.7b.

Macroscopic interstitial located in the lattice, true anti-singularity of the macroscopic vacancy

Now imagine that a cluster of interstitial type singularities, i.e. singularities carrying negative curvature charges, such as interstitial type prismatic dislocation loops, collapses on itself under the effect of attractive gravitational forces described in the previous chapter. If the initial cluster is neutral vis-à-vis the rotation charges, the individual singularities of the cluster combine, losing their own identity (dislocation or disclination loops) to form a single piece of macroscopic local lattice embedded within the lattice, and formed of N_I sites, as shown in figure 7.8.

This means that there is an excess of N_I lattice sites locally forming *a macroscopic interstitial*. The radius of this macroscopic embedding, assumed to be spherical, then is worth $R_{I0} = \sqrt[3]{3N_I/(4\pi n_0)}$ in the imaginary case where $\tau = 0$.

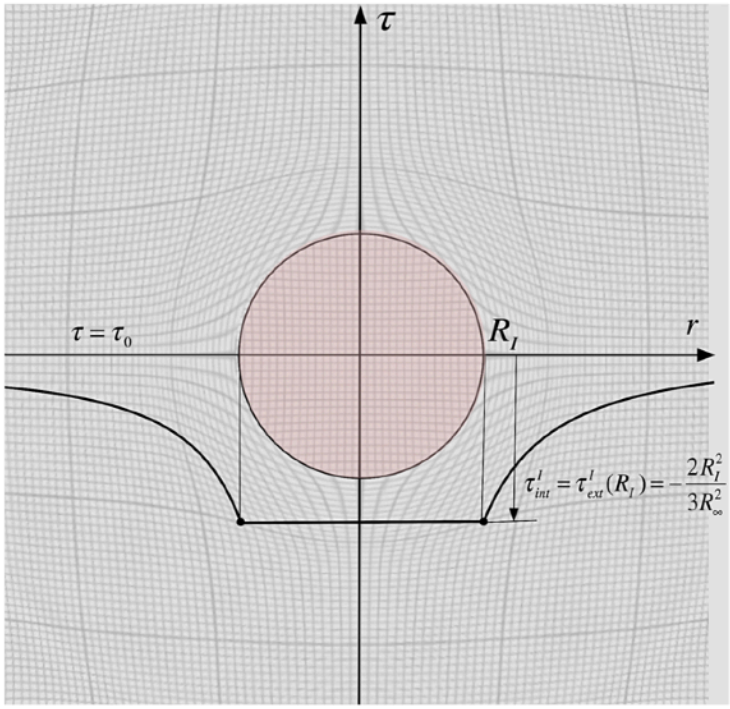
We can obviously consider that this macroscopic interstitial of N_I lattice sites corresponds in fact to *the anti-singularity of the macroscopic vacancy* of $N_L = N_I$ sites, in the sense that the combination of these two singularities completely restores the original lattice, since the N_L missing sites of the lattice are filled by the N_I interstitials.

In the presence of such an embedding of a piece of lattice within the lattice, there is obviously no coherence of the two lattices, and the condition of equilibrium amounts to the fact that the pressures at the interface are equal on the part and other of the interface, so that $p_{ext}(R_I) = p_{int}$, which actually comes down to that $\tau_{int}^{(I)} = \tau_{ext}^{(I)}(R_I)$. As the external field $\tau_{ext}^{(I)}(r)$ satisfies the divergence equation $\Delta\tau_{ext}^{(I)}(r) \cong 0$ drawn from Newton's first partial equation in figure 5.1, we deduce the solution for the external field, ensuring that the number of cells of lattices before introducing the singularity interstitial is equal to the number of cells of lattice after introduction of the singularity. Under this condition, one finds the fields of expansion perturbations outside and inside the macroscopic interstitial represented in figure 7.8.

Assuming that there exists in the vicinity of the macroscopic interstitial a field of expansion $\tau^{external}(R_I)$ due to the other singularities located in its vicinity, the real radius of the macroscopic interstitial will depend on the background expansion τ_0, on its own internal field of expansion internal and on the external field $\tau^{external}(R_I)$ by the relation also reported in figure 7.8. We can use this expression of R_I to express the field of external perturbations a little differently, in a form which shows that the external field of expansion perturbations is simply proportional to the number of additional sites agglutinated in the lattice.

The "gravitational" energy of the expansion field associated with this macroscopic interstitial is then written, taking into account the external and internal fields, as shown in figure 7.8.

Although *the macroscopic interstitial is the anti-singularity of the macroscopic vacancy,* its formation energy is always positive as long as $\tau_0 < K_0/2K_1$, and infinitely smaller than the formation energy of the macroscopic vacancy, which reveals *a colossal asymmetry between the two complementary macroscopic singularities*.



$$\begin{cases} \tau_{int}^{(I)} = \tau_{ext}^{(I)}(R_I) = -\dfrac{2R_I^2}{3R_\infty^2} << 1 \\ \tau_{ext}^{(I)}(r) = -\dfrac{2R_I^3}{3R_\infty^2}\dfrac{1}{r} = -\dfrac{a^3}{2\pi R_\infty^2} N_I \dfrac{1}{r} << 1 \\ R_I \cong \sqrt[3]{\dfrac{3N_I}{4\pi n_0}}\, e^{\frac{\tau_0 + \tau^{external}(R_I)}{3}} \end{cases}$$

$$E_{grav}^{(I)} \cong \frac{4\pi R_I^3}{3}\left(K_0 - 2K_1\tau_0\right) \cong \left(K_0 - 2K_1\tau_0\right)\frac{N_I}{n_0} e^{\tau_0 + \tau^{external}(R_I)}$$

Figure 7.8 *- the internal and external expansion fields of a macroscopic interstitial and the expressions of its radius and its gravitational energy*

In the presence of a macroscopic lattice interstitial, the first condition in figure 3.5 for a macroscopic interstitial to form a black hole implies that the critical radius of formation would be given by $r_{cr} = R_I^3 / (3R_\infty^2)$, and the second condition, namely that $r_{cr} > R_I$, implies that $R_I > \sqrt{3}R_\infty$. Obviously, we deduce that a macroscopic lattice interstitial, whatever its size and energy, *can never behave like a black hole*. This conclusion is extremely interesting insofar as we have here a topological object presenting a considerable asymmetry with its anti-singularity, the macroscopic vacancy, which necessarily becomes a black hole as soon as $\tau_0 > 1 - \tau^{external}(R_L)$.

About the possible analogy between vacancy singularities and black holes and between interstitial singularities and neutron stars

The macroscopic interstitial is the perfect anti-singularity of the macroscopic vacancy if

$N_I = N_L$ since the combination of the two singularities completely restores the perfect lattice. But there is *a colossal difference* between these two singularities, since their respective energies of formation are extremely different, and that the macroscopic vacancy becomes a black hole within the cosmological lattice as soon as $\tau_0 > 1$ while the macroscopic interstitial can never become black hole.

If we consider that these two topological objects can be formed by gravitational collapse of clusters of singularities, of a lacunar nature for the macro vacancy (for example lacunar prismatic dislocation loops) and of interstitial type for the macro interstitial (for example interstitial prismatic dislocation loops), we find by analogy the formation of black holes and neutron stars by gravitational collapse in the theory of gravitation. But if a sufficient initial mass of the cluster is a condition for arriving at a gravitational collapse, it would not be the initial mass of the cluster which conditions the evolution towards a black hole or towards a neutron star, but the very natureof the initial cluster.

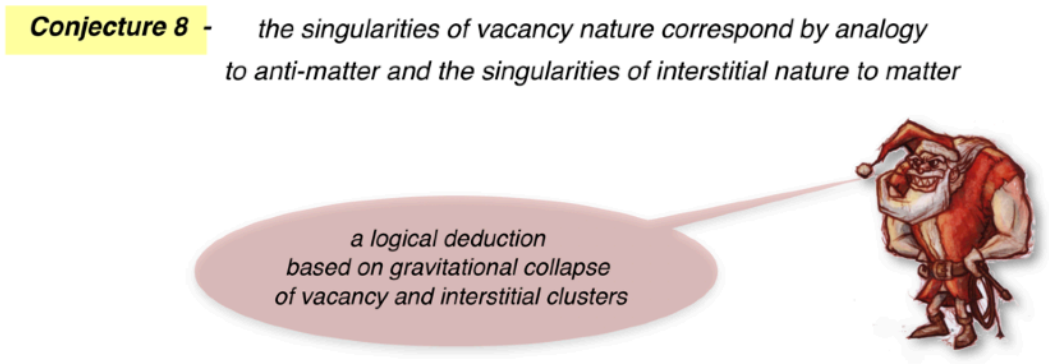


Figure 7.9 *- the eighth conjecture on the nature of matter and anti-matter*

If we then accept the conjecture 8 plotted in the figure 7.9, namely that the singularities of a vacancy nature correspond by analogy to anti-matter and the singularities of interstitial nature to matter, the black holes would then be residues of collapses of clusters of anti-matter of a lacunar nature and the neutron stars of the residues of collapsed clusters of matter of interstitial nature.

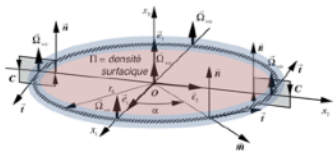
In this analogy, black holes, by virtue of their constitution of *"vacancy holes"* cannot keep any memory of the initial cluster of vacancy singularities from which they come, except *the quantity of microscopic vacancies*, that is to say the number of missing lattice sites in the initial vacancy singularities. On the other hand, neutron stars, by virtue of their constitution of *"interstitial embedding"* of pieces of lattice that are not coherent with the surrounding lattice, could conserve, apart from *the quantity of microscopic interstitials*, that is to say the number of excess lattice sites in the initial interstitial singularities, at least also the memory of the kinetic moment of rotation of the initial cluster of interstitial singularities from which they originate, in the form of *a very rapid rotation of the embedding,* which would correspond well with the enormous rotational speeds observed in the case of neutron stars, also called *pulsars* because of the electromagnetic pulses which they emit at fixed frequency following their very fast rotation.

Perturbations of the expansion associated with the twist disclination loop (BV)

The twist disclination loop has already been described in detail in figure 5.5. Concerning its "gravitational" properties, namely its external fields with long and short range of perturbations of the expansion, one can directly deduce them using the relations of figure 7.5, and report them in figure 7.10, using elsewhere the mass of inertia M_0^{BV} of the loop in place of its distortion energy E_{dist}^{BV} thanks to the expression of Einstein's relation. As these fields are perturbations of the volume expansion, they correspond in our analogy with the real world to *gravitational fields*, acting respectively at long distance *(LD)* and at short distance *(CD)*. In this form, the field of long distance expansion perturbations *(LD)* depends exclusively on the mass of inertia of the loop, and does not depend on the size of the loop, which further confirms our analogy between this field of expansionperturbations and a gravitational field.

The twist disclination loop (BV)

$$\begin{cases} \boldsymbol{q}_{\lambda BV} = -\pi R_{BV}^2 \Omega_{BV} = 2\pi R_{BV} \Lambda_{BV} = -\pi R_{BV} \vec{\boldsymbol{B}}_{BV} \vec{\boldsymbol{t}} \quad \textit{(charge of rotation)} \\ M_0^{BV} = \dfrac{2}{c_t^2}\left(K_2 + K_3\right)\zeta_{BV} R_{BV} \Lambda_{BV}^2 = \dfrac{1}{2c_t^2}\left(K_2 + K_3\right)\zeta_{BV} R_{BV}^3 \Omega_{BV}^2 \end{cases}$$



$$\begin{cases} \vec{\omega}_{ext}^{BV}(\vec{r}) = \dfrac{\boldsymbol{q}_{\lambda BV}}{4\pi} \dfrac{\vec{r}}{r^3} \qquad \textit{(divergent field of rotation)} \\ \tau_{ext\,LD}^{BV}(\vec{r}) \cong -\dfrac{c_t^2}{2\pi\left(K_0 - 2K_1(1+\tau_0) - 4K_2/3\right)R_\infty^2} \dfrac{M_0^{BV}}{r} \\ \tau_{ext\,CD}^{BV}(r) \cong \dfrac{c_t^2}{2\pi\left(K_0 - 2K_1(1+\tau_0) - 4K_2/3\right)} \dfrac{R_{BV} M_0^{BV}}{2\pi\zeta_{BV}} \dfrac{1}{r^4} \end{cases}$$

$$\begin{cases} E_{\boldsymbol{v}}^{BV} = \underbrace{\dfrac{1}{\gamma_t}\left(1 - \dfrac{\boldsymbol{v}^2}{2c_t^2}\right)E_{dist}^{BV}}_{E_{\boldsymbol{v}}^{dist}} + \underbrace{\dfrac{1}{\gamma_t}\dfrac{1}{2}M_0^{BV}\boldsymbol{v}^2}_{E_{\boldsymbol{v}}^{cin}} = \dfrac{E_{dist}^{BV}}{\gamma_t} = \dfrac{M_0^{BV} c_t^2}{\gamma_t} \\ E_{dist}^{BV} \cong M_0^{BV} c_t^2 \qquad \textit{if} \quad \boldsymbol{v} = 0 \\ E_{cin}^{BV} \cong \dfrac{1}{2} M_0^{BV} \boldsymbol{v}^2 \quad \textit{if} \quad \boldsymbol{v} \ll c_t \end{cases}$$

Figure 7.10 *- The essential properties of a twist disclination loop (BV)*

The expansion perturbation fields that we have just described obviously have an energy, which it would be desirable to compare with the elastic energy E_{dist}^{BV} of the screw loop. The total energy of the expansion perturbation field due to the elastic energy of the loop is deduced by the relation of figure 7.5 and is reported in figure 7.10. We can compare this energy with elastic energy, by relating it, and we get approximately $\boldsymbol{E}_{grav}^{(\boldsymbol{q}_{\lambda BV})} / \boldsymbol{E}_{dist}^{BV} \propto (K_1/K_0)(\vec{\boldsymbol{B}}_{BV}/R_{BV})^2 << 1$, because $K_1 << K_0$. It is therefore deduced that the "gravitational" energy $E_{grav}^{(\boldsymbol{q}_{\lambda BV})}$ of the twist disclination loop due to its elastic energy E_{dist}^{BV} is perfectly negligible compared to its elastic

energy in the perfect cosmological lattice.

In figure 7.10, we have reported all of the important properties that we have deduced so far for a twist disclination loop in a perfect cosmological lattice, namely its charge, its mass of inertia, its fields of rotation and disturbances of the long distance expansion, its elastic and kinetic energies, and finally its relativistic behavior. We see, among other things, that the mass of inertia M_0^{BV} of the loop does not only control the dynamic properties of the loop, like its kinetic energy E_{cin}^{BV}, but that it is also it which generates the gravitational fields $\tau_{ext\,LD}^{BV}(r)$ and $\tau_{ext\,CD}^{BV}(r)$ of external perturbations of expansion.

As the "gravitational" energy $E_{grav}^{(q_{\lambda BV})}$ of the expansion perturbation fields due to the rotation charge $q_{\lambda BV}$ of the twist disclination loop is perfectly negligible compared to the elastic energy E_{dist}^{BV} of the loop, this "gravitational" energy is not listed in the essential properties table.

As for the field of perturbations of the expansion within the torus itself surrounding the loop, and of the energy which is associated with this internal field, we will come back to this in detail later, when we will deal with the problem of spin and magnetic moment of the twist disclination loop.

Perturbations of the expansion associated with the edge prismatic loop (BC)

The prismatic edge dislocation loop we described in figure 5.6 has a curvature charge $q_{\theta BC}$, which makes it the elementary building block of *the lattice curvature charge* in our real world analogy. This charge is responsible for an external divergent flexion field, analogous to a geometric curvature field.

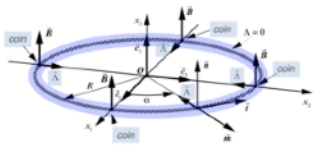
Knowing the elastic energy of the prismatic edge dislocation loop, we can use the relation of figure 7.1 to calculate the external field $\tau_{ext}^{(E)}(r)$ of expansion perturbations associated with the elastic energy E_{dist}^{BC} of the edge dislocation loop, here neglecting its energy potential V_{pot}^{BC}. One can also use the relation of figure 7.3 to calculate the external field $\tau_{ext}^{(q_{\theta BC})}(r)$ of expansion perturbations associated with the curvature charge $q_{\theta\,BC}$ of the loop. But using conjecture 5, that is $K_2 << K_3$, one can show that the first term $\tau_{ext}^{(E)}(r)$ is most likely far less than the second term $\tau_{ext}^{(q_{\theta BC})}(r)$. Thus, in the case of this loop, it is the expansion perturbation field due to the curvature charge which largely prevails over the expansion perturbation field due to the elastic distortion energy of the loop, unlike the twist disclination loop for which the energy of the perturbation field due to the distortion energy of the loop greatly outweighs the energy of the perturbation field due to the external field of rotation.

Apart from the mass of inertia M_0^{BC} of the edge loop, one can also introduce *an equivalent mass of curvature* $M_{curvature}^{BC}$ defined in figure 7.11, and which can be positive or negative depending on whether the edge loop is of a lacunar or interstitial nature. This mass of curvature then allows the gravitational field $\tau_{ext\,LD}^{BC}(r)$ to be written in the form shown in figure 7.11. According to the hypothesis of conjecture 5 ($K_2 << K_3$), the equivalent mass of curvature in the expression of the gravitational field satisfies the relationship $\left|M_{curvature}^{BC}\right| >> M_0^{BC}$. In addition, the equivalent mass of curvature can be positive or negative. These two facts will imply very surprising results as we will see later.

The expansion perturbation fields that we have just described obviously have an energy, which it would be desirable to compare with the elastic energy E_{dist}^{BC} of the edge loop. The

energy $E_{grav}^{(E)}$ of the expansion perturbation field due to the elastic energy of the loop is deduced thanks to the relations of figure 7.1. This energy can be compared with elastic energy, and it then comes the ratio $\boldsymbol{E}_{grav}^{(E)} / \boldsymbol{E}_{dist}^{BC} \propto (K_1 K_2^2 / K_0^3)(\vec{\boldsymbol{B}}_{BC} / R_{BC})^2 <<< 1$ as $K_1 K_2^2 <<< K_0^3$. We therefore deduce that the "gravitational" energy $E_{grav}^{(E)}$ of the edge dislocation loop due to its elastic energy E_{dist}^{BC} is perfectly negligible compared to its elastic energy in the perfect cosmological lattice.

Let's take a look at the energy $E_{grav}^{(\boldsymbol{q}_{\theta BC})}$ associated with the expansion perturbation field due to the curvature charge of the edge loop. Using the relation of figure 7.3, one obtains the energy of the field of "gravitation" associated with the charge of curvature, which can be compared with the elastic energy E_{dist}^{BC}, and one obtains the ratio $\boldsymbol{E}_{grav}^{(\boldsymbol{q}_{\theta BC})} / \boldsymbol{E}_{dist}^{BC} \propto K_1 / K_0 << 1$ as $K_1 << K_0$. We deduce again that the "gravitational" energy $E_{grav}^{(\boldsymbol{q}_{\theta BC})}$ of the dislocation edge loop due to its curvature charge $\boldsymbol{q}_{\theta BC}$ is perfectly negligible compared to its elastic energy E_{dist}^{BC} in the perfect cosmological lattice.



The prismatic edge dislocation loop (BC)

$$\boldsymbol{q}_{\theta BC} = -2\pi\vec{\boldsymbol{n}}\left(\vec{\boldsymbol{t}} \wedge \vec{\Lambda}_{BC}\right) = 2\pi\vec{\Lambda}_{BC}\vec{\boldsymbol{m}} = -2\pi\vec{\boldsymbol{n}}\vec{\boldsymbol{B}}_{BC} \begin{cases} >0 & \text{if vacancy loop} \quad (\vec{\boldsymbol{n}}\vec{\boldsymbol{B}}_{BC} < 0) \\ <0 & \text{if interstitial loop} \quad (\vec{\boldsymbol{n}}\vec{\boldsymbol{B}}_{BC} > 0) \end{cases}$$

$$M_0^{BC} = \frac{E_{dist}^{BC}}{c_t^2} \cong \left(\frac{K_2}{K_3}\right)^2 \frac{K_3}{c_t^2}\zeta_{BC} R_{BC} \vec{\Lambda}_{BC}^2 = \left(\frac{K_2}{K_3}\right)^2 \frac{K_3}{c_t^2}\zeta_{BC} R_{BC} \vec{\boldsymbol{B}}_{BC}^2$$

$$M_{curvature}^{BC} = -\frac{2\pi K_2}{5c_t^2} R_{BC}^2 \vec{\boldsymbol{n}}\vec{\boldsymbol{B}}_{BC} = \frac{K_2}{5c_t^2} R_{BC}^2 \boldsymbol{q}_{\theta BC} \begin{cases} >0 & \text{if vacancy loop} \quad (\boldsymbol{q}_{\theta BC} > 0) \\ <0 & \text{if interstitial loop} \quad (\boldsymbol{q}_{\theta BC} < 0) \end{cases}$$

$$\left|M_{curvature}^{BC}\right| >> M_0^{BC}$$

$$\vec{\chi}_{ext}^{BC}(\vec{r}) = \frac{\boldsymbol{q}_{\theta BC}}{4\pi} \frac{\vec{r}}{r^3}$$

$$\tau_{ext}^{BC}(r) \cong -\frac{c_t^2}{2\pi\left(K_0 - 4K_2/3 - 2K_1(1+\tau_0)\right)R_\infty^2} \frac{M_{curvature}^{BC} + M_0^{BC}}{r}$$

$$E_{\boldsymbol{v}}^{BC} = \underbrace{\frac{1}{\gamma_t}\left(1 - \frac{\boldsymbol{v}^2}{2c_t^2}\right)E_{dist}^{BC}}_{E_{\boldsymbol{v}}^{dist}} + \underbrace{\frac{1}{\gamma_t}\frac{1}{2}M_0^{BC}\boldsymbol{v}^2}_{E_{\boldsymbol{v}}^{cin}} = \frac{E_{dist}^{BC}}{\gamma_t} = \frac{M_{0\epsilon}^{BC} c_t^2}{\gamma_t}$$

$$E_{dist}^{BC} \cong M_0^{BC} c_t^2 \quad \text{if} \quad \boldsymbol{v} = 0$$

$$E_{cin}^{BC} \cong \frac{1}{2} M_0^{BC} \boldsymbol{v}^2 \quad \text{if} \quad \boldsymbol{v} << c_t$$

Figure 7.11 *- The essential properties of a prismatic edge dislocation loop (BC)*

In figure 7.11, we have reported all the important properties that we have deduced so far for an edge dislocation loop in a perfect cosmological lattice, namely its curvature charge, its mass of inertia, its equivalent gravitational mass, its fields of flexion and of perturbations of the expansion of the long distance, its elastic and kinetic energies, and finally its relativistic behavior. One notes there that it is its elastic energy E_{dist}^{BC}, and as a consequence its mass of inertia M_0^{BC} which is deduced from it, which controls its dynamic properties, like its kinetic

energy E_{cin}^{BC} , but that it is its equivalent gravitational mass $M_{curvature}^{BC}$ of curvature which essentially controls its external gravitational field $\tau_{ext\,LD}^{BC}(r)$ of expansion perturbations at long-distance.

As the "gravitational" energies $E_{grav}^{(E)}$ and $E_{grav}^{(q_{\theta BC})}$ of the fields of expansion perturbations due to the elastic energy E_{dist}^{BC} and to the curvature charge $q_{\theta BC}$ of the dislocation edge loop are perfectly negligible compared to the elastic energy E_{dist}^{BC} of the loop, these "gravitational" energies do not appear in the table of essential properties.

On the other hand, the equivalent gravitational mass of curvature $M_{curvature}^{BC}$ is not only much greater than the mass of inertia M_0^{BC}, but it can even be negative in the case of edge loops of interstitial nature. This result turns out to be very surprising since it corresponds to the possible existence of a negative gravitational field, and it is in fact extremely promising by the novelty which it brings in our approach compared to the General Relativity of Einstein. The possible consequences of this astonishing result will be explored in detail in the rest of the book.

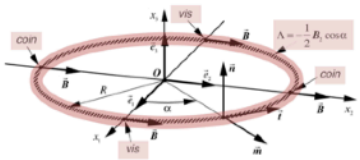
Perturbations of the expansion associated with the mixed slip dislocation loop (BM)

The mixed slip dislocation loop that we have described in figure 5.7 has neither rotation charge nor curvature charge, but on the other hand a dipolar moment of rotation field $\vec{\omega}_{dipolar}^{BM}(r,\theta,\varphi)$, analogous to an electric dipolar moment. Knowing the elastic energy of distortion of the mixed slip dislocation loop in a perfect cosmological lattice, we can use the relation of figure 7.1 to calculate the external field $\tau_{ext}^{E}(r)$ of expansion perturbations associated with this elastic energy E_{dist}^{BM}, by neglecting here the potential energy V_{pot}^{BM} . On the other hand, it also becomes interesting here to replace in the expression $\tau_{ext}^{E}(r)$ the distortion energy E_{dist}^{BM} by the mass of inertia M_0^{BM} of the loop using the relation $E_{dist}^{BM} = M_0^{BM} c_t^2$. *It comes the expression of* $\tau_{ext}^{E}(r)$ *reported in figure 7.12.*

The slip mixed dislocation loop (BM)

$$\begin{cases} M_0^{BM} \cong \dfrac{K_2 + K_3}{4c_t^2} \zeta_{BM} R_{BM} \vec{B}_{BM}^2 \end{cases}$$

$$\begin{cases} \exists \quad \text{an external dipolar rotation field} \quad \vec{\omega}_{dipolar}^{BM}(r,\theta,\varphi) \\ \tau_{ext}^{BM}(r) \cong -\dfrac{c_t^2}{2\pi\left(K_0 - 4K_2/3 - 2K_1(1+\tau_0)\right)R_\infty^2} \dfrac{M_0^{BM}}{r} \end{cases}$$



$$\begin{cases} E_{\boldsymbol{v}}^{BM} = \underbrace{\dfrac{1}{\gamma_t}\left(1 - \dfrac{\boldsymbol{v}^2}{2c_t^2}\right)E_{dist}^{BM}}_{E_{\boldsymbol{v}}^{dist}} + \underbrace{\dfrac{1}{\gamma_t}\dfrac{1}{2}M_0^{BM}\boldsymbol{v}^2}_{E_{\boldsymbol{v}}^{cin}} = \dfrac{E_{dist}^{BM}}{\gamma_t} = \dfrac{M_0^{BM}c_t^2}{\gamma_t} \\ E_{dist}^{BM} \cong M_0^{BM}c_t^2 \qquad si \quad \boldsymbol{v} = 0 \\ E_{cin}^{BM} \cong \dfrac{1}{2}M_0^{BM}\boldsymbol{v}^2 \qquad si \quad \boldsymbol{v} << c_t \end{cases}$$

Figure 7.12 *- The essential properties of a slip mixed dislocation loop (BM)*

The expansion perturbation fields we have just found obviously have an energy, which it would be desirable to compare with the elastic energy E_{dist}^{BM} of the loop. The energy $\boldsymbol{E}_{grav}^{(E)}$ of the expansion perturbation field due to the elastic energy of the loop is deduced thanks to the relations of figure 7.1, and can be compared directly with the elastic energy E_{dist}^{BM} to derive the ratio $\boldsymbol{E}_{grav}^{(E)} / \boldsymbol{E}_{dist}^{BM} \propto (K_1 / K_0)(\vec{\boldsymbol{B}}_{BM} / R_{BM})^2 << 1$ as $K_1 << K_0$. We therefore deduce that the "gravitational" energy $E_{grav}^{(E)}$ of the mixed dislocation loop due to its elastic energy E_{dist}^{BM} is perfectly negligible compared to its elastic energy in the perfect cosmological lattice.

In figure 7.12, we have plotted the set of properties of a mixed slip dislocation loop in the perfect cosmological lattice. If this loop has neither rotation charge nor curvature charge generating external fields of rotation or flexion at long range, it is provided with an external field of dipolar rotation $\vec{\omega}_{dipolar}^{BM}(r,\theta,\varphi)$, at short range, analogous to an electric dipolar field. The external field of expansion perturbations is due to the elastic distortion energy of the loop, and depends on its mass of inertia M_0^{BV}.

On the other hand, since the "gravitational" energy $E_{grav}^{(E)}$ of the expansion perturbation fields due to the elastic energy E_{dist}^{BM} of the mixed dislocation loop is perfectly negligible compared to the elastic energy E_{dist}^{BM} of the loop, this "gravitational" energy is not included in the table of essential properties.

About the various properties of the elementary topological loops

We have demonstrated above all the essential properties of the three elementary loops that can be encountered in the perfect cosmological lattice, and we have notably established the expressions of the fields of perturbations of external expansion, which correspond to "gravitational fields" of expansion of the loops, and which will play a very important role thereafter.

We have shown that the energies $E_{grav}^{(E)}, E_{grav}^{(\boldsymbol{q}_{\lambda BV})}, E_{grav}^{(\boldsymbol{q}_{\theta BC})}$ of the "gravitational" fields associated with the elastic energies, the rotation charges and the curvature charges of the elementary loops are perfectly negligible with respect to the elastic energies associated with these loops, and that we can therefore perfectly ignore them in our calculations.

One can wonder if the elementary topological loops can be black holes. To do this, we apply *the conditions of the black holes of figure 3.5* to the expansion fields $\tau_{ext\,LD}^{BV}(r)$, $\tau_{ext}^{BC}(r)$ and $\tau_{ext}^{BM}(r)$ respectively. By applying these conditions to *the twist disclination loop* in the case where $\tau_0 << \tau_{0cr}$, we obtain the following condition $\left|\vec{\boldsymbol{B}}_{BV}\right| \geq 2R_\infty (2\pi / \zeta_{BV})^{1/2}$ so that this loop is a black hole, which would imply that the pseudo-vector of Burgers of the loop is of the order of magnitude of the radius of the lattice, which obviously makes no sense.

The same goes for *the mixed slip dislocation loop*, since the condition for it to be a black hole becomes $\left|\vec{\boldsymbol{B}}_{BM}\right| > 4R_\infty (2\pi / \zeta_{BM})^{1/2}$. In the case of *the prismatic edge dislocation loop*, the condition is expressed differently since the gravitational mass is worth $M_{curvature}^{BC} + M_0^{BC}$. Like $\left|M_{curvature}^{BC}\right| >> M_0^{BC}$, we get $R_{BC}\left|\vec{\boldsymbol{B}}_{BC}\right| \geq 10 K_0 R_\infty^2 / K_2$. But since the module K_2 must be much smaller than the module K_0 according to conjecture 6, this condition can obviously never be satisfied. Consequently, it is clear that *the three elementary loops forming the basis of the microscopic topological singularities cannot be black holes in the domain* $\tau_0 << \tau_{0cr}$.

In the analogy between our approach and the major theories of physics, the three types of elementary loops that we have discussed in this chapter have a number of amazing and

remarkable properties that we will list here:

- they are respectively the elementary bricks of *the electric charge*, *the curvature charge* and *the electric dipolar moment*, from which it might be possible to form dispirations, by more or less complicated combination of several loops, in order to find topological singularities which are analogues of the elementary particles of our real world,
- their rest energy and their kinetic energy are essentially *confined in the toric field surrounding the loops*,
- as the energy associated with the "gravitational" fields of expansion perturbations is perfectly negligible, *they perfectly satisfy Einstein's relation*, which is surprisingly obtained in our approach as a purely classical property of topological singularities within the lattice, without appeal to a principle of relativity,
- they perfectly satisfy *special relativity*, with a completely original explanation of relativistic energy $E_{\boldsymbol{v}}^{loop} = E_{\boldsymbol{v}}^{dist} + E_{\boldsymbol{v}}^{cin} = E_{dist}^{loop} / \gamma_t = M_0^{loop} c_t^2 / \gamma_t$ as the sum of a relativistic term of elastic distortion energy and a relativistic term of kinetic energy,
- they satisfy *a relativistic dynamic equation* reported in figure 6.7,
- the twist disclination loop, carrying a rotation load similar to the electric charge, satisfies *Maxwell's equations* and *Lorentz force*,
- the three types of loops present a long distance field of perturbations of the volume expansion, which is the analog of *a gravitational field decreasing in* $1/r$ and which *depends only on a gravitational mass* $M_0^{loop} + M_{curvature}^{loop}$ of the loops composed of the mass of inertia and of the equivalent curvature mass of the loop, without directly depending on the size R_{loop} of the loops,
- the gravitational masses of the twist disclination loop and the mixed slip dislocation loop are strictly equal to their masses of inertia, while the gravitational mass of the prismatic edge dislocation loop consists of the mass of inertia and the mass of curvature of the loop, with a mass of curvature much higher than the mass of inertia, and which may even be negative in the case of loops of interstitial nature,
- the gravitational mass of the prismatic edge dislocation loop contains two terms: the first dominant term of mass of curvature $M_{curvature}^{BC}$ is positive or negative depending on whether the loop is of lacunar or interstitial type, and the second term of mass of inertia M_0^{BC} is always positive. This means that the gravitational mass is not symmetrical between a loop of vacancy nature and a loop of interstitial nature. There appears here *a weak asymmetry in the absolute value of the gravitational mass between an interstitial loop and its lacunar anti-loop*, which is expressed by the following three relationships:

$$M_{curvature}^{BC(l)} + M_0^{BC(l)} > 0 \quad ; \quad M_{curvature}^{BC(i)} + M_0^{BC(i)} < 0 \quad ; \quad \left| M_{curvature}^{BC(i)} + M_0^{BC(i)} \right| \tilde{<} M_{curvature}^{BC(l)} + M_0^{BC(l)}$$

- all these properties are perfectly analogous to the fundamental properties of elementary particles in the real world, except for the gravitational mass of the prismatic edge dislocation loops, loops which have in fact a strong analogy with neutrinos. This very special property of prismatic edge dislocation loops will be discussed in *the following chapters, in which we will discuss the gravitational interaction between loops.*