Chapter 8

Newtonian gravitation and general relativity

A detailed study of the gravitational interactions intervening between loops of twist disclination leads to a strong analogy with the gravitation of Newton, but presenting however some differences as regards the gravitational attraction at short range and especially the dependence of the constant of gravitation in the volume expansion of the lattice.

We are then interested in *the Maxwellian formulation* of the spatio-temporal evolution equations presented previously in chapter 4, which corresponded to the expression of local laws of physics, such as electromagnetism, as seen by the imaginary external observer ***GO*** *(Great Observer).* We show in this chapter that it is possible to imagine a local observer ***HS*** *(Homo Sapiens)* which would be intimately linked to the local coordinate system, because it itself consists of a cluster of topological singularities of the lattice. This observer can only know local rulers and clocks in his local coordinate system, which are influenced by the local expansion field so that they make Maxwell's equations invariant with respect to volume expansions. A relativistic notion of time then appears for the local observers ***HS***, for whom the speed of the transverse waves is measured as a universal constant, whereas this depends very strongly on the local volume expansion if it is measured by the observer ***GO*** outside the lattice.

The gravitational interactions thus obtained have very strong analogies with *Newton's Gravitation* and with *Einstein's General Relativity*. Perfectly analogous points are discussed in detail, such as the perfect analogy with *Schwarzschild's metric* at a great distance from a massive object and *the curvature of the wave rays* by a massive object.

But we also show that our Eulerian approach of the cosmic lattice brings new elements to the theory of Gravitation, in particular modifications of the Schwarzschild metric at very short distance from a singularity, and a better understanding of the critical rays associated with black holes: *the radii of the perturbation sphere and of the point of no return are both similar and equal to the Schwarzschild radius*, while *the limit radius for which the dilation of the observer's time would tend towards infinity becomes zero*, so our approach is not limited to the description of a black hole beyond the Schwarzschild sphere.

While the clusters of singularities composed of loops of twist disclination satisfy Newton's gravitation, and most of the results of General Relativity, we will finally ask ourselves the question of how to behave, towards the gravitation, other topological singularities, such as edge dislocation loops, mixed dislocation loops, macroscopic vacancies and macroscopic interstitials. We will thus deduce all the long-range gravitational interaction forces between the various topological singularities and their behaviors.

Newtonian gravitational interaction of clusters of twist disclination loops

In the previous chapter, it was shown that the mass of inertia of twist disclintion loops is

extremely higher than the mass of inertia of edge and mixed dislocation loops, so that perturbations of the expansion field will be mainly caused by the twist disclination loops. This is why we will start by analyzing the long-range expansion perturbation fields due to them.

For a screw loop, the mass of inertia depends on the square of the rotation charge $\boldsymbol{q}_{\lambda BV}$ (see figure 7.10). If such a loop is in an expansion field τ, the dependence of the mass of inertia M_0^{BV} and the charge of rotation $\boldsymbol{q}_{\lambda BV}$ of this loop is related to the dependence $c_t = c_{t0}\,\mathrm{e}^{\tau/2}$ and the dependence of R_{BV} and Ω_{BV} in the local expansion. One could think a priori that the radius of the loop is linked to the pitch of the lattice, so that $R_{BV} = R_{BV0}\,\mathrm{e}^{\tau/3}$, and that the angle of rotation Ω_{BV} must correspond to an angle satisfying *the symmetry of the lattice*, namely for example a multiple of $\pi/2$ for a cubic lattice, or a multiple of $\pi/3$ for a hexagonal lattice or for a face-centered cubic lattice, so that Ω_{BV} should not depend on the background expansion of the lattice.

However, we do not know at all the exact nature of the cosmological lattice, so that, for the sake of generality, we will assume *a priori* an unknown dependence of R_{BV} and Ω_{BV} on the background expansion of the lattice, as if the loop could undergo an extension of its radius and / or a torsion dependent on the local volume expansion of the lattice, writing by assumption that $R_{BV} = R_{BV0}\,\mathrm{e}^{\alpha_{BV}\tau}$ and $\Omega_{BV} = \Omega_{BV0}\,\mathrm{e}^{\beta_{BV}\tau}$, where α_{BV} and β_{BV} are constants which will have to be determined subsequently, as indicated in figure 8.1. On the basis of this assumption, one can deduce the dependencies in the lattice expansion τ of the charge $\boldsymbol{q}_{\lambda BV}$ and the mass of inertia $\boldsymbol{M}_0^{BV}$ of a loop, then the dependencies in τ of the charge $\boldsymbol{Q}_\lambda^{cluster}$, the mass of inertia $\boldsymbol{M}_0^{cluster}$ and the energy $\boldsymbol{E}^{cluster}$ of a cluster of rotation charges.

Two clusters of loops located at a distance d from each other will interact with each other via their "gravitational" fields of long-range expansion perturbations (figure 8.1). We can calculate the energies of the clusters (1) and (2) from their mass of inertia, by introducing into the energy the total value of the expansion at the heart of each cluster. Indeed, the two clusters distant from the distance d are respectively immersed in the expansion perturbation field of the other cluster. As we know that the elastic energy of a loop is essentially located in the vicinity close to the loop, one is assured that the elastic energy of the cluster is essentially located in the core of the cluster, so that their respective energy is influenced by the presence of the other cluster and so that there appears an increase ΔE_{grav} in the energy of the two interacting clusters, which depends on the distance d separating the two clusters. The total interaction force between the two clusters is then given by the derivative with respect to d of the energy variation ΔE_{grav} of the two clusters, which leads to an interaction force shown in figure 8.1. This expression shows the product of the gravitational masses divided by the squared distance separating the two clusters, which furiously recalls a term from Newton's law of gravitation. Let's introduce *a "gravitational constant"* $\boldsymbol{G}_{grav}$ function of the lattice constants α_{BV} and β_{BV} and of lattice elastic constants, but also of the local volume expansion τ of the lattice.

The total force $\boldsymbol{F}_{grav}(d)$ of gravitational interaction between the two clusters takes a form that can be recognized when we develop the exponential terms it contains, since it shows *the expression of Newton's gravitational force* between the two loops, corrected by additional terms in the bracket, which are the results of second-order developments, so that they are certainly smaller than 1. For rather low mass densities of the clusters and large distances d between the clusters, the second order terms in the bracket can be neglected, so that *a perfect analogy with*

Newton's gravitational interaction $\boldsymbol{F}_{grav}(d) \cong \boldsymbol{G}_{grav} M_{0(1)}^{cluster} M_{0(2)}^{cluster} / d^2$ of the real world is found, and $\boldsymbol{G}_{grav}$ becomes the "gravitational constant", which must be very small since it occurs R_{∞}^2 in the denominator.

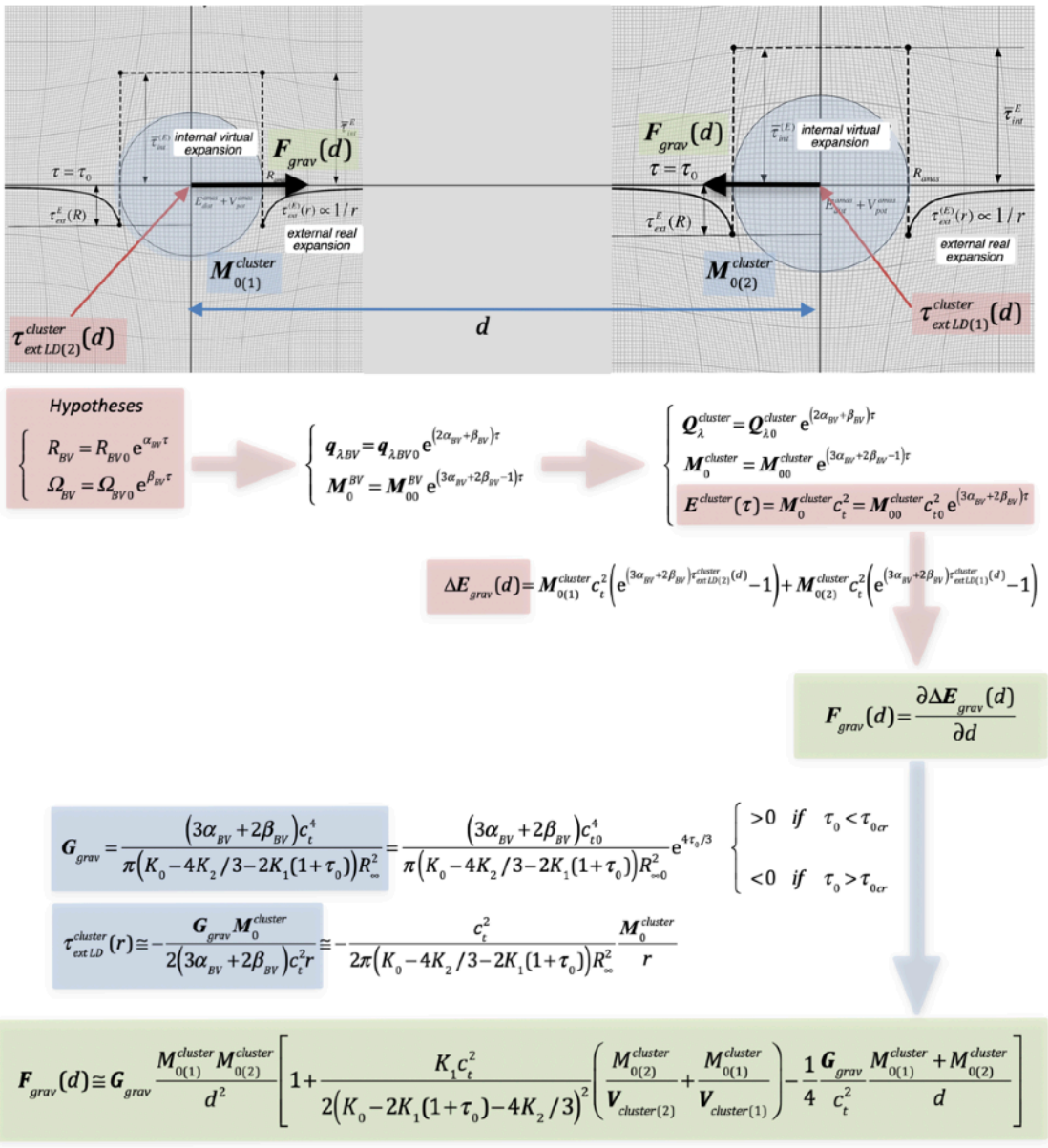

Figure 8.1 *- The «gravitational» interaction force between two clusters of singularities with rotation charge*

Let us then try to understand what the second-order terms actually represent in the full expression $\boldsymbol{F}_{grav}(d)$ of Figure 8.1. The term dependent on cluster mass densities $M_{0(i)}^{cluster} / \boldsymbol{V}_{cluster(i)}$ is a correction to the "gravitational constant» $\boldsymbol{G}_{grav}$ independent of the distance d between clusters, which actually comes from the non-linear behaviour of the function $\overline{\tau}_{int}^{(E)}(x)$ shown in Fig. 7.0 when the cluster mass density exceeds a certain critical value. As for the term dependent on d , it brings a correction to the interaction force when the distance between clusters becomes small. This term comes from the development of the

exponential terms of the energy variation ΔE_{grav} of the two clusters.

What first appears when looking at the expression of the "constant" of gravitation is that it is in fact in no way a constant since it depends on the average state of expansion τ_0 of the lattice via the values of c_t^4 and R_∞^2, as well as by the value of τ_0 intervening in the denominator in factor of the module K_1, which makes the value of $\boldsymbol{G}_{grav}$ positive if $\tau_0 < \tau_{0cr}$ and negative if $\tau_0 > \tau_{0cr}$. This strong dependence of $\boldsymbol{G}_{grav}$ on the background expansion of the lattice should undoubtedly play a key role in the evolution of the Universe during its cosmological expansion. We will come back to this subject later.

At a shorter distance d, there appears *a correction to Newton's law* expressed by a multiplicative term. This second-order term in the expression of the approximate law of gravitation will modify the interactions between two clusters when they are very close to each other. But unlike the results obtained in General Gravitation using the Schwarzschild metric, which predicts a small increase in the force of attraction at very short distances, the second-order corrective term in our approach leads to a small decrease in $1/d^3$ of the attraction force at very short distances.

For example, in the case of the planet Mercury which is quite close to the Sun, the calculation of the corrective term provides the value $\boldsymbol{G}_{grav} M_{Sun} / (4c_t^2 d_{Mercury}) \cong 5{,}6 \cdot 10^{-7}$ with $M_{Sun} \cong 2 \cdot 10^{30} \left[kg \right], d_{Mercury} \cong 5{,}8 \cdot 10^{10} \left[m \right], \mathrm{G}_{grav} \cong 6{,}6 \cdot 10^{-11} \left[m^3 / kg \cdot s^2 \right], c_t^2 \cong 10^{17} \left[m^2 / s^2 \right]$ which gives for the period of revolution of Mercury 88 days ($7{,}6 \cdot 10^6$ seconds), an increase of 2.128 seconds compared to the value calculated with Newton's law of gravitation.

The local rulers and clock of an HS observer

Consider a local reference frame $\boldsymbol{O}x_1x_2x_3$ defined by the ***GO*** observer *(the imaginary Great Observer)* from its absolute frame of reference $\boldsymbol{Q}\xi_1\xi_2\xi_3$. This local coordinate system $\boldsymbol{O}x_1x_2x_3$ is in fact a convenience used by ***GO*** to solve the problems of local evolution of the solid lattice, in particular in regions of the solid having a non-zero volume expansion, but which can be considered as constant and homogeneous in the vicinity of the origin of the coordinate system $\boldsymbol{O}x_1x_2x_3$, for example using *the Maxwellian formulation* described in chapter 4.

But now imagine that there is indeed another category of local observer that we called observer ***HS*** *(Homo Sapiens)* and which is really found in the local coordinate system $\boldsymbol{O}x_1x_2x_3$ because it is itself constituted from the topological singularities of the lattice, and in particular from elementary loops of twist disclination which interact with each other via their rotation field generated by their own rotation charge. In his local frame, the ***HS*** obviously does not have access to the global view of the lattice in the absolute reference frame, since he only knows the conditions of the local lattice in which he lives. The ***HS*** observer can therefore only define his own rulers $\vec{e}_{yi}$ in his coordinate system $\boldsymbol{O}y_1y_2y_3$, by defining them from the linear dimensions of the objects contained in the lattice in which he lives. This is illustrated in figure 8.1 for two observers ***HS*** and ***HS'*** living in two different locations of the lattice, where the respective volume expansions of the lattice τ and τ' are different.

If the lattice locally exhibits a certain volume expansion τ, the ***HS*** rules should satisfy a relationship of the type $\vec{e}_{yi} = \mathrm{e}^{a\tau} \vec{e}_i$ where the constant a is not defined a priori and must therefore be determined. This implies that the rulers of an ***HS*** will be of different length than those of the ***GO*** if the lattice is locally shrinking ($\tau < 0$) or expanding ($\tau > 0$). If a certain point

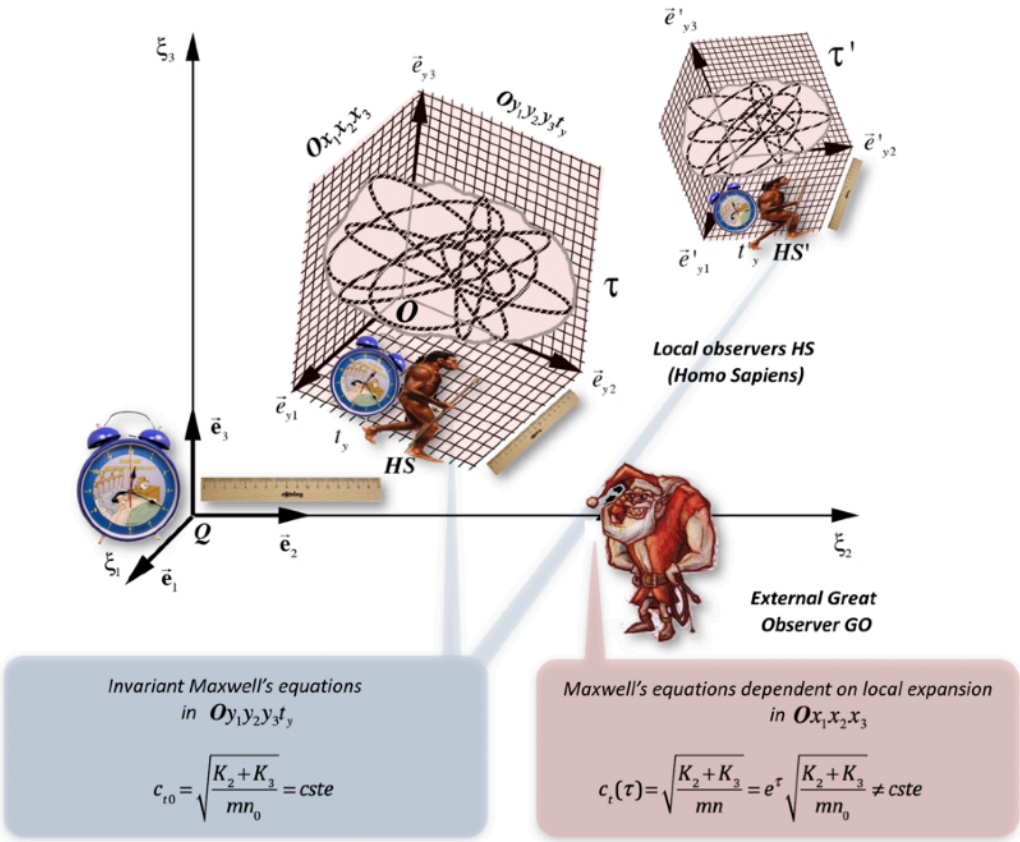

Figure 8.2 *- Local rulers and clocks in the reference frames of* ***HS*** *and* ***HS'***

in space, identified by the vector $\vec{r}$, is observed simultaneously by the ***GO*** and by ***HS***, the vector will be written $\vec{r}=\Sigma\, x_i\vec{e}_i=\Sigma\, x_i\,\mathrm{e}^{-a\tau}\,\vec{e}_{yi}=\Sigma\, y_i\vec{e}_{yi}$ in $Ox_1x_2x_3$ and in $Oy_1y_2y_3$ respectively so that the coordinates of the point of the space is transformed by relationship $y_i=\mathrm{e}^{-a\tau}\,x_i$.

On the other hand, the time measured in $Ox_1x_2x_3$ at $\tau\neq 0$ must also be different from the time measured when $\tau=0$, so that ***HS*** own clock in its frame $Oy_1y_2y_3$ must indicate a time t_y different from the absolute time t of the ***GO***, but linked to it by the relationship $t_y=\mathrm{e}^{b\tau}\,t$. Regarding the ***HS*** clock, he will have to build it locally since it does not have access to the absolute time of the ***GO***. It will then be logical for him to build a simplistic clock using one of his local rulers and the speed of the transverse waves which he can measure in the coordinate system of his own frame of reference. Consider a ruler of length rule d_0 measured by ***GO*** in a lattice with zero volume expansion. The length of the same ruler placed in the ***HS*** frame and measured by the ***GO*** becomes $d=d_0\,\mathrm{e}^{a\tau}$. To traverse the distance of this ruler, the transverse waves use a period of time Δt measured by the ***GO***, and given by $\Delta t=\Delta t_0\,\mathrm{e}^{-b\tau}$. In the coordinate system $Ox_1x_2x_3$, placed in a volume expansion region τ other than zero, the initial length d_0 of the ruler therefore becomes equal to $d_0\,\mathrm{e}^{a\tau}$, and the travel time of the ruler becomes $\Delta t_0\,\mathrm{e}^{-b\tau}$, so that the speed of the transverse waves in the presence of the expansion τ is equal to $c_t=d/\Delta t=c_{t0}\,\mathrm{e}^{(a+b)\tau}=c_{t0}\,\mathrm{e}^{\tau/2}$. This implies that time will flow differently for ***HS***

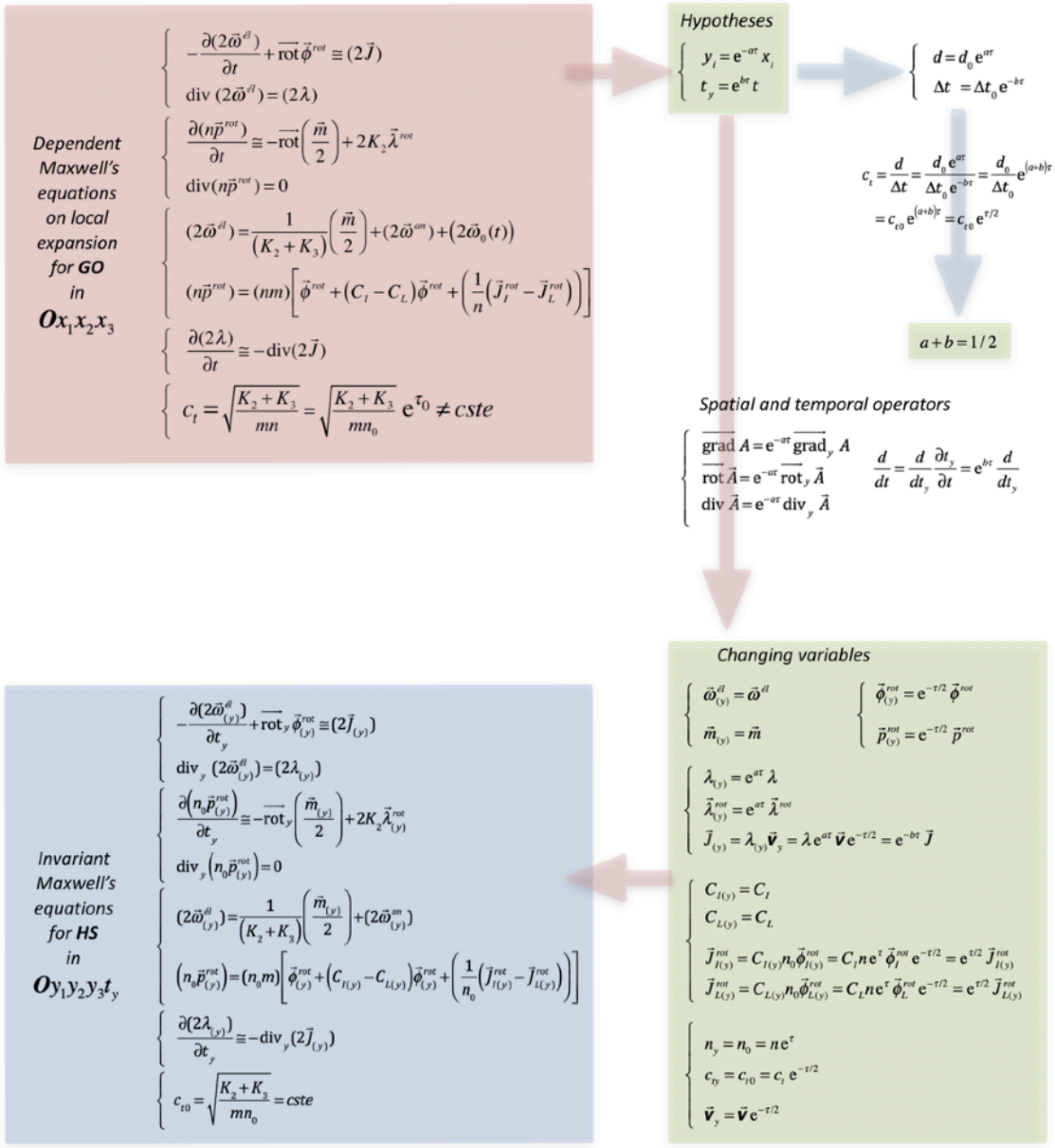

Figure 8.3 *- The transition from the local frame $Ox_1x_2x_3$ of **GO** to the local frame $Oy_1y_2y_3$ of **HS***

in $Oy_1y_2y_3$ if the lattice is contracting or expanding. We then have the following relationship between the values of a and b due to the operation of the ***HS*** clock: $a+b=1/2$.

In figure 8.3, one deduces all the expressions making it possible to pass from the local reference frame $Ox_1x_2x_3$ of the ***GO*** to the local reference frame $Oy_1y_2y_3$ of the ***HS***. On the basis of the assumptions for the rulers and the clock of ***HS***, one expresses the formulas of passage of the operators of time and space, and one applies them to the equations of Maxwell expressed in $Ox_1x_2x_3$ to obtain these same equations in $Oy_1y_2y_3$. By making a change of variables in the fields of rotation, torque, velocity and momentum, as well as on the quantities associated with the densities and fluxes of rotation charges and the quantities associated with the concentrations of vacancies and of interstitial and to the fluxes of these defects, we obtain the Maxwell's equations as formulated by ***HS*** in his local frame of reference $Oy_1y_2y_3$.

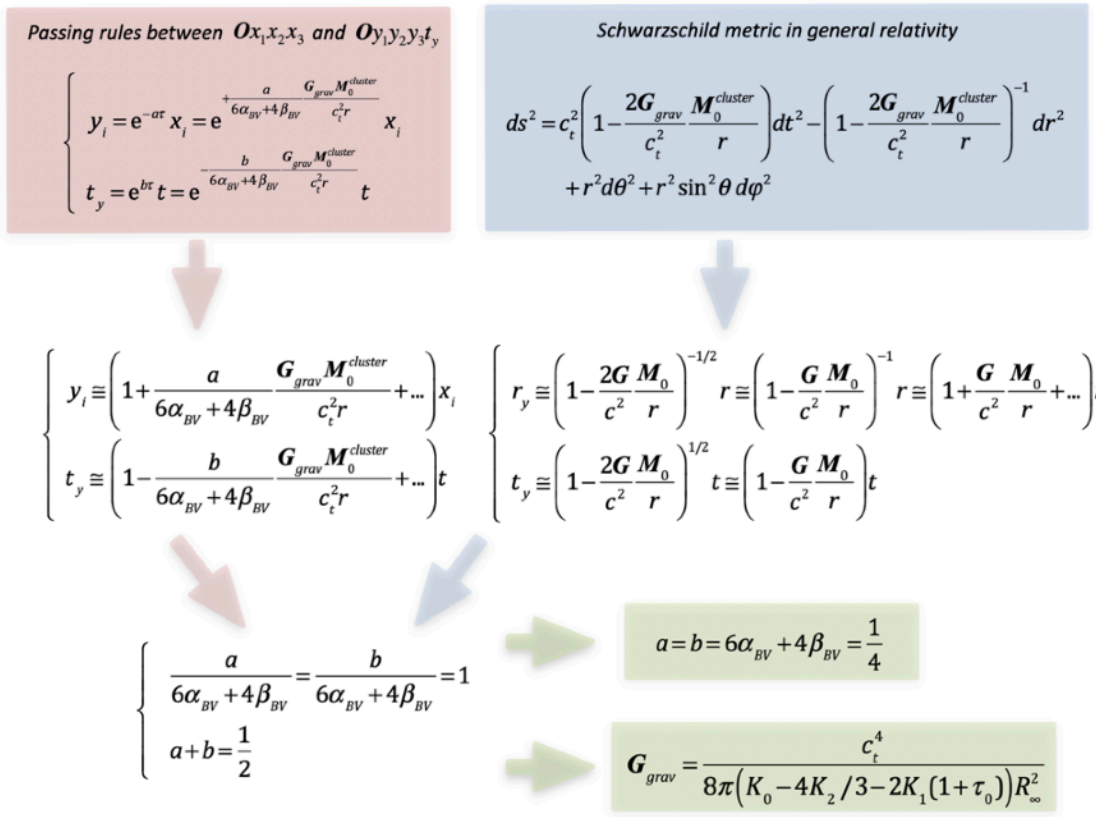

Figure 8.4 *- Comparison of passing rules between the local frame $Ox_1x_2x_3$ of **GO** and the local frame $Oy_1y_2y_3$ of **HS** with the transformation rules experimentally verified of the Schwarzschild metric in general relativity*

In these expressions of change of variables, it is perfectly logical to have posed $\vec{\omega}_{(y)}^{él} = \vec{\omega}^{él}$ and $\vec{m}_{(y)} = \vec{m}$ since these quantities are associated with measurements of angles of rotation, and to have chosen the relations $\vec{\phi}_{(y)}^{rot} = e^{-\tau/2} \vec{\phi}^{rot}$ and $\vec{p}_{(y)}^{rot} = e^{-\tau/2} \vec{p}^{rot}$, since these quantities are in fact associated with velocities. By this change in the variables, we then obtain *a set of invariant Maxwell equations in the frame $Oy_1y_2y_3t_y$ of the observer **HS*** as illustrated in figures 8.2 and 8.3, i.e. equations which no longer depend on the local expansion τ. We note, among other things, that the speed of the transverse waves becomes an invariant constant for the ***HS*** observer, whatever the state of expansion of the lattice in which he lives. We also note that only the quantities associated with the densities and the flow of rotation charges are transformed in a dependent way in the parameter a, while all the other quantities are transformed in a logical and predictable way.

The fact that the Maxwell equations of the ***HS*** observers are invariant (independent of local expansion τ) implies that the local ***HS*** observers are perfectly incapable of measuring the local state of expansion of the lattice in which they live on the sole basis of electromagnetic measurements based on Maxwell's equations. In particular, the measurement of the speed of the transverse waves by the ***HS*** observers always provides an invariant value, whatever the local state of expansion of the lattice. Thus, local observers ***HS*** are essentially subject to the

physical laws corresponding to electromagnetism, and are only aware of the gravitational effects associated with the expansion field through indirect observations of their effects, such as the movement of the planets or the slowing down of their clocks in a gravitational field. This is why the ***HS*** will have to seek to explain the phenomena related to gravitation by ad-hoc theories (Newton's Gravitation, General Relativity) which seem a priori independent of the laws of electromagnetism, but which they will obviously seek to unify.

Figure 8.2 illustrates well here the existence of a strong analogy between our approach and *Einstein's General Relativity Theory*. Indeed, the rulers and the clock of an ***HS*** living in a certain place of the lattice depend on the local volume expansion τ of the lattice, in a way analogous to what stipulates the General Relativity for the rulers and the clock of an observer located in a given gravitational field. In the case illustrated in figure 8.2, we understand that the lattice plays the role of an *"ether"* which imposes the size of the rulers of the ***HS*** observer, while it is the celerity of the transverse waves within the lattice (in fact the speed of information transport) which imposes the speed of the ***HS*** clock.

On the other hand, the existence of three degrees of freedom on the parameters a, b, α_{BV} and β_{BV} is quite surprising, because it implies that there is still a possible choice at this level, which cannot be determined on the basis of our knowledge of the cosmological lattice. In fact, an arbitrary choice of "free" values of a, α_{BV} and β_{BV} should not lead to inconsistencies in the system, and the resulting cosmological lattice could be perfectly viable. Consequently, we can consider that the parameters a, α_{BV} and β_{BV} are really properties specific to the cosmological lattice, in fact intrinsic constants of the lattice, in the same way as are the elastic modules K_i or the mass m of inertia per cell of the lattice. The determination of these constants therefore necessarily involves experimentation, in other words by measuring the real properties of the cosmological lattice associated with these two constants.

By combining the various relations obtained in figure 8.3, we can write the transformation rules to pass from the local reference frame $\boldsymbol{O}x_1x_2x_3$ of the ***GO*** into the local reference frame $\boldsymbol{O}y_1y_2y_3t_y$ of an ***HS*** placed at a distance r in the gravity field of a cluster of mass M_0^{amas}, according to unknown parameters a, α_{BV} and β_{BV}, as shown in figure 8.4. These transformation equations can, at a certain sufficient distance r from a cluster of mass $M_0^{cluster}$, be written approximately by developpement of the exponentials.

In General Relativity, the dependencies of the radial ruler and the clock of an observer subjected to the gravitational field of an object of mass M_0^{amas} are deduced from *the Schwarzschild metric*. This is obtained in the case of a massive object with spherical symmetry by postulating an invariant metric with respect to the rotations, which is also shown in figure 8.4. From this Schwarzschild metric, we deduce that the radial ruler and the observer's clock depend symmetrically on the distance r separating him from the object of mass M_0, with expressions similar to *the simplified expressions obtained in our approach*, but with a coefficient 1 in front of $\boldsymbol{G}_{grav}M_0^{cluster}/c_t^2r$. However, the dilation of time in a gravity field, represented by the second relation deduced from the metric of Schwarzschild, has



Karl Schwarzschild (1873-1916)

been verified experimentally with very great precision[1], even on elevation differences as small as 1 meter on the surface of the earth, and this effect is currently taken into account in very precise navigation systems, such as *GPS*. This experimentally verified effect can therefore be used to adjust the determination of the parameters a, α_{BV} and β_{BV}, by ensuring that the long-distance relationships of our approach correspond to the relationships deduced from the Schwarzschild metric. By making this comparison, we get that $a = b = 6\alpha_{BV} + 4\beta_{BV} = 1/4$.

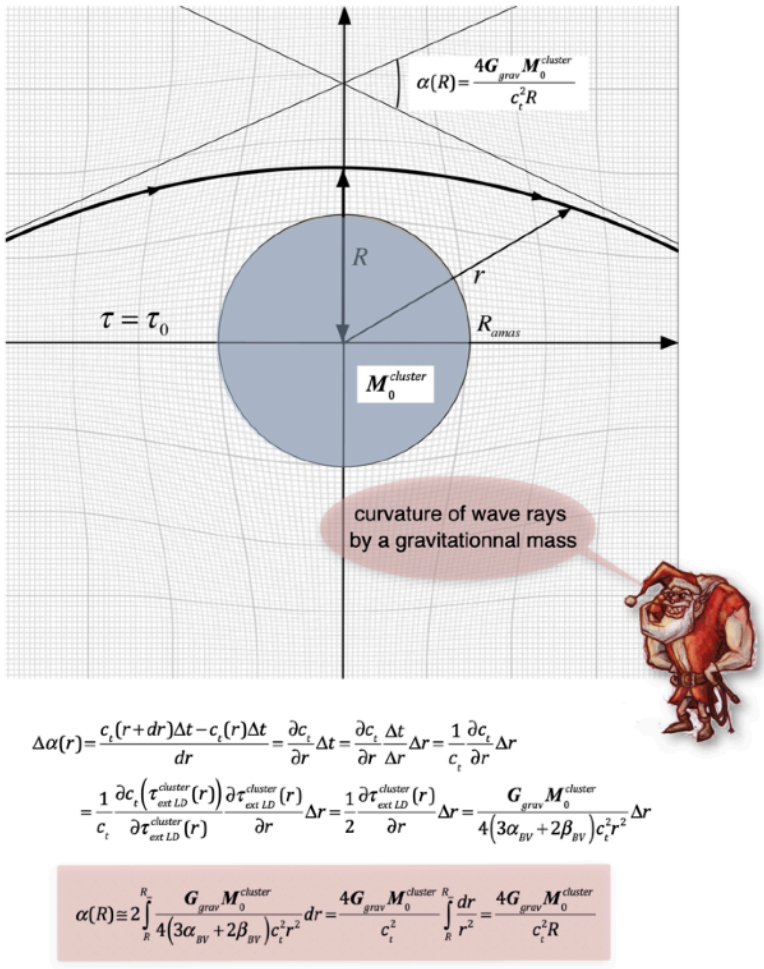

Figure 8.5 *- curvature of transversal wave rays in the vicinity of a massive cluster*

We can still verify the analogy of our approach with Einstein's General Relativity, for example by calculating the curvature of the rays of transverse waves in the vicinity of a massive cluster, as illustrated in figure 8.5, since the measurement of this effect at the beginning of the 20th century was the first experimental verification of Einstein's General Relativity.

In the vicinity of a massive cluster, the celerity $c_t(r)$ of the transverse waves depends on the distance r to the center of the cluster. This implies that the directions of two radii of waves perpendicular to a straight line passing through the center of the cluster, and passing to the

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[1] see for example: *C. W. Chou, D. B. Hume, T. Rosenband, D. J. Wineland : "Optical Clocks and Relativity", Science, vol. 329, 5999, pp. 1630-1633*

distances r and $r+\Delta r$ from the cluster will have an infinitesimal angle $\Delta\alpha(r)$. We can use the dependence of the celerity $c_t(r)$ in the gravitational field $\tau_{ext\,LD}^{cluster}(r)$ to express $\Delta\alpha(r)$. As the wave will travel the path shown in figure 8.5, the tangents to infinity of the incident wave and of the deflected wave form a total angle $\alpha(R)$ depending on the minimum distance R of the wave from the center of the cluster. Half of the total angle $\alpha(R)$ can then be deduced approximately by integrating $\Delta\alpha(r)$ for distances r to the center of the cluster ranging from R to R_∞, and the desired value of the deflection angle $\alpha(R)$ is obtained. It turns out that this value is exactly the same as that obtained in general relativity, which is logical since we used Schwarzschild's metric to calibrate the constants of our approach. General Relativity provides for a curvature worth $\alpha(R)\cong 4\boldsymbol{G}_{grav}M_0^{cluster}/c_t^2R$, and the calculation of this value in the case of a light ray with a grazing incidence with the sun gives a deflection angle of 1,75'' of arc.

As for the experimental values of the deflection of light by the sun, measured by Eddington at the beginning of the 20th century (May 1919) during an eclipse of the sun, they gave approximately 1,98''$\pm$0,12'' of arc (at Sobral in Brazil) and 1,61''$\pm$0.31" of arc (in Sao Tomé-et-Principe in the Gulf of Guinea), which corresponds fairly well to the value calculated theoretically by General Relativity, despite the many difficulties of the experimental measurements.

During the Lorentz transformation described in chapter 6, the transformation laws are symmetrical with respect to time and space, just like the transformation laws in General Relativity in the case of the Schwarzschild metric. Among all the possible results that one could have imagined for the parameters a, b, α_{BV} and β_{BV}, the results obtained from the time dilation and the curvature of the wave rays in weak gravity field conveniently provide symmetric transformation laws, since we get that $a=b$. We can therefore issue here a conjecture 9, which combines the effects of time dilation and curvature of the wave rays in a weak gravity field, and reported in figure 8.6

Conjecture 9 - *the metric of our theory in a weak gravitational field must be the same as Schwarzschild's metric in General Relativity*

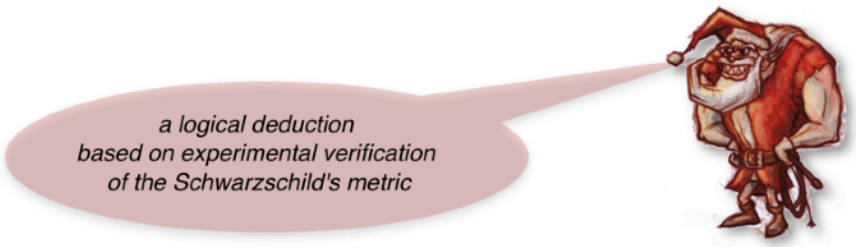

Figure 8.6 *- The ninth conjecture on the long distance metric of a massive object*

We still have a degree of freedom to choose the values of the parameters α_{BV} and β_{BV}, which must be linked by expression $3\alpha_{BV}+2\beta_{BV}=1/8$. However, knowing exactly the values of α_{BV} and β_{BV} is not important with regard to the gravitational properties of the loop, since the fact that $3\alpha_{BV}+2\beta_{BV}=1/8$ implies that the mass of inertia M_0^{BV} and the distortion energy E^{BV} of a loop, as well than the quantities $\boldsymbol{G}_{grav}$, $\tau_{ext\,LD}^{cluster}(r)$ and $\boldsymbol{F}_{grav}(d)$ do not explicitly depend on α_{BV} and β_{BV}.

The exact choice of parameters α_{BV} and β_{BV} obviously depends on the deep nature of the cosmological lattice, which is beyond the scope of this treaty. But the simplest solution to imagine would be that β_{BV} is equal to zero, because, in this case, the torsion Ω_{BV} of the twist disclination loop would become a constant independent of the expansion of the lattice, for example a multiple of $\pi/2$ for a cubic lattice, or a multiple of $\pi/3$ for a hexagonal lattice or for a face-centered cubic lattice, which obviously seems the most logical.

But a zero value of β_{BV} would also imply that α_{BV} would be worth 1/24, and therefore that the radius $R_{BV} = R_{BV0}\,\mathrm{e}^{\tau/24}$ of the loop would depend much less on the expansion of the lattice than the pitch $a = a_0\,\mathrm{e}^{\tau/3}$ of the lattice or the length $\vec{e}_{yi} = \mathrm{e}^{\tau/4}\,\vec{e}_i$ of the ***HS*** observer ruler.

Unable to make an explicit safe choice of the values of the parameters α_{BV} and β_{BV}, we will leave this problem open, while remembering that there is this degree of freedom for the choice of their values, which must simply be linked by the expression $3\alpha_{BV} + 2\beta_{BV} = 1/8$.

Analogies and differences from general relativity

Let us assume that a twist disclination loop of inertia mass M_0^{BV} is at the confines of the lattice, and that a cluster of singularities with an inertia mass $M_0^{cluster} >> M_0^{BV}$ is at the center of the lattice. The gravitational force acting on the loop is therefore written approximately $\boldsymbol{F}_{grav}^{BV}(r) \cong \boldsymbol{G}_{grav} \boldsymbol{M}_0^{cluster} \boldsymbol{M}_0^{BV} / r^2$. Under the effect of this force, the twist disclination loop will undergo an acceleration in the direction of the cluster of singularities, and the potential gravitational energy of the loop will gradually transform into kinetic energy of the loop which will be worth $\boldsymbol{E}_{kinetic}^{BV}(r) \cong \boldsymbol{G}_{grav} \boldsymbol{M}_0^{cluster} \boldsymbol{M}_0^{BV} / d$, depending from its distance d from the cluster. But the kinetic energy of the loop is written $\boldsymbol{E}_{kinetic}^{BV} = \boldsymbol{M}_0^{BV} \boldsymbol{v}^2 / 2$ in the non-relativistic case so that the speed of the loop as a function of the distance d separating it from the cluster depends on *the gravitational constant* and *the mass of inertia of the cluster:* $\boldsymbol{v}^2(R) = 2\boldsymbol{G}_{grav} \boldsymbol{M}_0^{cluster} / d$. As *the mass of inertia of the loop does not intervene in this relation*, this relation remains valid even for relativistic speeds of the loop. However, we know that the total relativistic energy $\boldsymbol{v}^2(R) = 2\boldsymbol{G}_{grav} \boldsymbol{M}_0^{cluster} / d$ of the loop will tend towards infinity when its speed $\boldsymbol{v}(R)$ will tend towards the celerity c_t of the transverse waves, so that the following condition $\boldsymbol{v}^2(R) \le c_t^2$ must be satisfied before reaching the speed limit. This condition implies the existence of *a critical distance* $d_{cr} = 2\boldsymbol{G}_{grav} \boldsymbol{M}_0^{cluster} / c_t^2$ for which the energy of the loop becomes infinite. This critical distance d_{cr} only depends on the mass $M_0^{cluster}$ of the cluster, and obviously only exists if the radius of the cluster is less than d_{cr}. It is called *the Schwarzschild radius* of the cluster and in fact corresponds to the limit beyond which the loop cannot irreversibly leave the cluster since it would then require infinite energy to do so. Thus, the mass $M_0^{cluster}$ of the cluster whose radius $R_{cluster}$ would satisfy the condition $R_{cluster} < d_{cr}$ would in fact be *a black hole* which would irreparably absorb any singularity which would have the misfortune to approach it at a distance $R \le R_{Schwarzschild}$. This *Schwarzschild radius*[2] $R_{Schwarzschild} = 2\boldsymbol{G}_{grav} \boldsymbol{M}_0^{cluster} / c_t^2$ is obtained in exactly the same way in Relativity, so that it is identical in our approach and in General Relativity.

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[2] *Schwarzschild, K. (1916). On the gravitational field of a point mass according to Einstein's theory. Sitzber. Preuss. Akad. Wiss., Physik-Math. Kl., Vol. 189, pp.189-196. (translated by Helga and Roger Stuewer).*

In our approach, we have already approached the notion of black hole by defining in figure 3.5 the conditions for a singularity of the gravitational field to behave like a black hole *with respect to transverse waves*, by defining the radius r_{cr} of the sphere of perturbations around the singularity, namely the sphere beyond which any transverse wave is trapped by the singularity. By then applying these conditions to the cluster of mass $M_0^{cluster}$ generating the gravitational field $\tau_{ext\,LD}^{cluster}(r)$ of figure 8.1, one obtains the expression of *the radius* r_{sphere} *of the sphere of perturbations* of the cluster in the form $r_{sphere} = 2\boldsymbol{G}_{grav}\boldsymbol{M}_0^{cluster}/c_t^2$, i.e. the same value than the radius of the Schwarzschild sphere. In General Relativity, we define *the sphere of photons*, i.e. the limit in the vicinity of a black hole from which no photon can escape from the black hole, whose radius is equal to $R_{photons\,sphere} = 3\boldsymbol{G}M/c^2$, namely 3/2 of the Schwarzschild radius. We will come back to this disturbing difference with our approach.

In our approach, the proper time of an ***HS*** is actually written exactly as an exponential term proportional to $-1/r$. With this exponential expression, we see that the proper time of the observer ***HS*** which approaches at very short distance from a cluster *expands infinitely when the distance tends towards zero*. In General Relativity, it is said that the proper time of the ***HS*** seems to expand infinitely when the ***HS*** approaches a critical limit distance r_{cr} calculated on the basis of the Schwarzschild metric, when $t_y \to 0$, which leads to the critical distance $r_{cr} = \boldsymbol{G}_{grav}\boldsymbol{M}_0^{cluster}/c_t^2$. This limit distance is therefore smaller than the radius of the Schwarzschild sphere, the point of no return for a black hole. Now, it seems rather difficult to imagine that the time of an ***HS*** seems to stop when the ***HS*** reaches this critical distance, and it is surprising that this critical distance is half the radius of Schwarzschild, and not simply the radius of Schwarzschild itself, or else the zero radius as in our approach.

Perfect cosmological lattice		*General Relativity*
$\begin{cases} y_i = e^{\frac{G_{grav}M_0^{cluster}}{c_t^2 r}} x_i \\ t_y = e^{\frac{G_{grav}M_0^{cluster}}{c_t^2 r}} t \end{cases}$	*Transformation laws*	$\begin{cases} y_i = \left(1 + \frac{G}{c^2}\frac{M_0}{r} + \dots\right) x_i \\ t_y = \left(1 - \frac{G}{c^2}\frac{M_0}{r}\right) t \end{cases}$
$R_{Schwarzschild} = 2G_{grav}M_0^{cluster}/c_t^2$	*Schwarzschild radius*	$R_{Schwarzschild} = 2G_{grav}M_0^{cluster}/c_t^2$
$R_{photon\,sphere} = 2G_{grav}M_0^{cluster}/c_t^2$	*Photon sphere radius*	$R_{photon\,sphere} = 3G_{grav}M_0^{cluster}/c_t^2$
$R_{time\,dilation \to \infty} \to 0$	*Radius of infinite dilation of time*	$R_{time\,dilation \to \infty} \cong G_{grav}M_0^{cluster}/c_t^2$
Curvature vector by flexion		*Einstein's curvature tensor*
$\vec{\chi} = -\frac{4K_2/3+2K_1}{2(K_2+K_3)}\overrightarrow{\text{grad}}\tau + \frac{1}{2(K_2+K_3)}\left(n\frac{d\vec{p}}{dt} - \overrightarrow{\text{grad}}F^{def}\right) + \frac{K_3}{K_2+K_3}\vec{\lambda}$		$G = 8\pi T$ *Energie-momentum tensor*
«Gravitational» field	*Energy-momentum vector*	*Flexion charge*
Divergence of Newton's equation: equation of motion for "gravitational" distortions		*Divergence of the energy-momentum tensor: gravitational motion equation*
$\vec{\nabla}\vec{\chi} = \text{div}\left[\frac{1}{2(K_2+K_3)}\left(n\frac{d\vec{p}}{dt} - \overrightarrow{\text{grad}}F^{def}\right) - \frac{4K_2/3+2K_1}{2(K_2+K_3)}\overrightarrow{\text{grad}}\tau + \frac{K_3}{K_2+K_3}\vec{\lambda}\right] = \theta$		$\vec{\nabla} \cdot \boldsymbol{T} = 0$

Figure 8.7 *- Comparative behavior of the perfect cosmological lattice with the general relativity of Einstein*

In General Relativity, the simplest black holes are characterized by three critical radii: *the radius of the photon sphere* which is worth $R_{photons\ sphere} = 3\boldsymbol{G}M/c^2$, *the radius of the horizon or point of no return*, also called *the Schwarzschild radius*, which is given by $R_{Schwrzschild} = 2\boldsymbol{G}M/c^2$, and *the radius for which the dilation of the observer's time tends to infinity*, which is approximately equal to $R_{time\ dilation \to \infty} \cong \boldsymbol{G}M/c^2$. The fact that there are three different radii for black holes in General Relativity is quite intriguing, as is the existence of a nonzero radius for which the dilation of the observer's time tends to infinity. It is mainly for this reason that it is said that it is not possible to describe by General Relativity the physics of objects that fall into a black hole beyond the Schwarzschild sphere.

In our approach, on the contrary, the radii of the perturbation sphere and of the point of no return are both similar to the Schwarzschild radius ($2\boldsymbol{G}M/c^2$), which is very satisfying for the mind, since it thus only exists a single boundary representing *the horizon of a black hole*. On the other hand, there is no limit radius in our approach for which the dilation of the observer's time would tend towards infinity, so that our approach is not limited to the description of a black hole beyond the Schwarzschild sphere. Now our approach is equivalent to General Relativity as long as the gravitational field is weak and satisfies the condition $\tau_{ext\ LD}^{cluster}(r) < 1$. The reason is that these are the two experimentally verified effects, namely the time dilation and the curvature of wave rays in weak fields which were chosen to make our approach identical to that of General Relativity in weak gravitational field. On the other hand, our approach becomes different for strong gravitational fields, as shown by the compared expressions of the transformation laws in figure 8.7, which explains the differences that we have just described with regard to the different characteristic radii of black holes.

About the formal analogy between the 3D spatial curvature equation of the cosmological lattice and the Einstein's equation of the 4D field of curvature in general relativity

The local spatial curvature of the cosmological lattice, as measured by the ***GO*** observer, is characterized by *the curvature vector by flexion* $\vec{\chi} = -\overrightarrow{\mathrm{rot}}\,\vec{\omega} + \vec{\lambda}$, perfectly described in figure 2.6. It is remarkable to note that this flexion vector can be directly obtained from *the Newton's equation of the cosmological lattice* (figure 3.1). This gives the value of the flexion field in the cosmological lattice in the form shown in figure 8.7. We therefore deduce that the existence of a local topological curvature of the lattice seen by the ***GO*** via the flexion vector depends on three terms at the same time:

- the gradient of the local volume expansion, which is nothing other than the gradient *of the "gravitational" field* τ within the lattice,
- the temporal variations of the local momentum of the lattice and the gradient of the density of elastic energy $F^{él}$ stored in the lattice, term which one could qualify as *"energy-momentum vector"* due to the singularities present in the lattice,
- the density $\vec{\lambda}$ of flexion charges within the lattice, which reflects the presence of *topological singularities within the lattice*, such as dislocations and / or disclinations.

On the other hand, for a local observer ***HS***, both *its rulers and its clock depend on the local volume expansion*, so that an equation similar to the flexion field equation should necessarily

become *a four-dimensional equation of curvature of space-time*, which we are obviously not going to try to establish here.

The operation of *taking the divergence of the curvature field*, i.e. $\mathrm{div}\,\vec{\chi} = \mathrm{div}\,\vec{\lambda} = \theta$, allows writing a second equation shown in figure 8.7. This relationship shows that *the divergence of the flexion vector is equal to the density of curvature charges* due to the topological singularities contained in the lattice, and is therefore zero if there are no curvature charges. Moreover, in the case where there are no curvature charges, *the divergence of the bending vector is nothing other than the divergence of the Newton's equation of the lattice*, i.e. the equation of motion for the divergent part of the distortions, namely the “gravitational” distortions by volume expansion.

The first relation giving the spatial curvature of the lattice from the Newton's equation of the lattice is the three-dimensional analog of *the Einstein's four-dimensional field equation of General Relativity*[3], which is written $\boldsymbol{G} = 8\pi\boldsymbol{T}$, in which $\boldsymbol{G}$ is *the famous curvature tensor of Einstein (Einstein tensor),* which is expressed in terms of *the Ricci curvature tensor* $\boldsymbol{G}_{\mu\nu} = R_{\mu\nu} - g_{\mu\nu}R/2$, corresponding to a certain part of *the Riemann tensor* which describes the curvatures of space-time. As for the tensor $\boldsymbol{T}$, it is *a "geometric object"* called *the stress-energy tensor*, or *the energy-momentum tensor,* or *the stress-energy-momentum tensor,* which characterizes the matter contained in space.

This Einstein field equation shows how the energy-momentum tensor of matter generates an average curvature of space-time in its vicinity. Among other things, it allows you to calculate the static curvature field of a massive object, or the generation of gravitational waves by a massive mobile object. In addition, it also contains the equations of motion ("Newton's equations") for matter whose energy-momentum tensor generates the curvature of space-time.

In the case of Einstein's field equation, it should also be noted that the energy-momentum tensor is a tensor with zero divergence $\vec{\nabla} \bullet \boldsymbol{T} = 0$, which ensures that the laws of conservation of energy and angular momentum are respected . This equation $\vec{\nabla} \bullet \boldsymbol{T} = 0$ therefore represents in fact *the equation of motion of matter* in General Relativity.

There is very clearly a strong analogy between Einstein's field equation $\boldsymbol{G} = 8\pi\boldsymbol{T}$ and the equation giving the flexion field $\vec{\chi}$ in the case of the cosmological lattice, described in table 8.7, because the latter also connects a “geometric vector” of spatial curvature to a kind of “energy-momentum vector” within the solid lattice, which contains at the same time the temporal variations of the local momentum of the lattice, the gradient of the local volume expansion, and the gradient of the density $F^{él}$ of elastic energy stored in the lattice, quantities which are all influenced by the presence of torsion charges or flexion charges within the lattice. In addition, this equation for the field $\vec{\chi}$ of curvature by flexion derives directly from the Newton's equation of the lattice. However, unlike Einstein's field equation, which describes the curvatures of 4-dimensional space-time, this equation is deduced by the ***GO*** which is lucky to have an absolute clock, so that *there is no "curvature of time" for him*, and that, consequently, his equation of curvature is purely spatial in 3 dimensions. On the other hand, if we considered the way that "homo sapiens" ***HS*** observers should have to describe the gravitational behavior they observe, *they should take into account the "curvature of time"* since their clocks depend on the local state of expansion of the lattice where they are.

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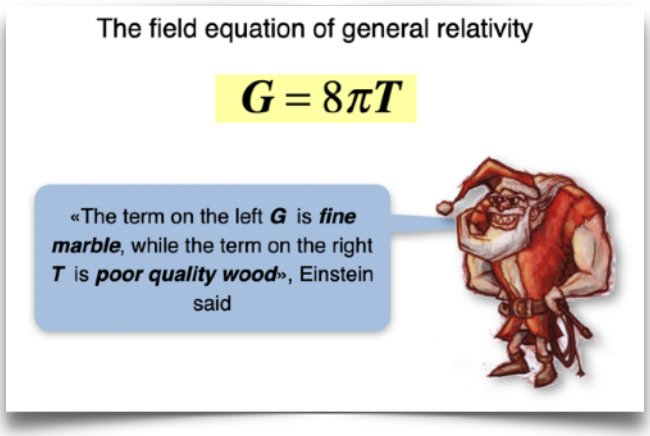
[3] *a very goof account of Einstein's General Relativity can be found in the book:*
Ch. W. Misner, K. S. Thorne, J. A. Wheeler, ***Gravitation****, W.H. Freeman and Co, San Francisco 1973*

In the case of Einstein's field equation, the equation of the divergence $\vec{\nabla} \bullet \boldsymbol{T} = 0$ of the energy-momentum tensor represents *the equation of gravitational movements of matter* in General Relativity, just like the equation $\vec{\nabla}\,\vec{\chi} = \mathrm{div}\,\vec{\chi} = \mathrm{div}\,\vec{\lambda} = \theta$ of the divergence of the curvature represents the divergence of the Newton's equation of motion of the lattice, that is to say the equation of motion for the "gravitational" distortions of the lattice by volume expansion.

But the most important difference between Einstein's field equations and cosmological lattice equations is that *the concept of curvature charges associated with topological singularities does not exist in Einstein's equations*, which appears in the fact that the divergence of the energy-momentum tensor is always zero, so that there is no equation similar to the equation $\mathrm{div}\,\vec{\chi} = \mathrm{div}\,\vec{\lambda} = \theta$ in General Relativity. We will see later in this book that the concept of curvature charges, absent in modern theories of relativity, quantum physics and particle physics, can explain a large number of facts misunderstood by modern theories of physics.



It is very tempting to say here that *the concept of curvature charges*, and that *the pure geometric equation* $\mathrm{div}\,\vec{\chi} = \mathrm{div}\,\vec{\lambda} = \theta$ answers the question that Einstein asked when he said that, in the field equation $\boldsymbol{G} = 8\pi\boldsymbol{T}$, the term on the left, the curvature tensor of Ricci $\boldsymbol{G}$, is made *of fine marble*, while the right term, the energy-momentum tensor $\boldsymbol{T}$, is made *of poor quality wood,* meaning that the right term is only a phenomenological representation of the matter injected into the field equation, which is not derived directly from a prime principle as is the term on the left. In our approach, on the other hand, the right term of the equation for the flexion field is made *of fine marble*, since it is derived directly from a prime principle which is the Newton's equation of the lattice.

About the dependence of the topological singularities on the expansion of the lattice

In the case of a twist disclination loop ***(BV)***, we have seen that the dependences in volume expansion of the radius of the loop, its angle of twist, its rotational charge and its mass of inertia can be written by introducing the constants α_{BV} and β_{BV}, which makes it possible to deduce several important relationships for the twist disclination loop in the presence of a background expansion field τ_0 and / or an external expansion field $\tau^{external}$ generated by other singularities , relations which are reported in the table of figure 8.8.

In the case of a prismatic edge dislocation loop ***(BC)***, we do not know a priori the dependencies in volume expansion of the radius of the loop, the Burgers vector and the charge of curvature. We therefore introduce new expansion constants α_{BC} and β_{BC} in such a way that we can write the relations for the radius R_{BC} of the loop, for the Burgers vector $\vec{\boldsymbol{B}}_{BC}$ and for the curvature charge $\boldsymbol{q}_{\theta BC}$ in the form shown in table 8.8. These expansion constants α_{BC} and β_{BC} are not known. The dependence of the radius of the loop should certainly be similar to that

The twist disclination loop **(BV)**, source of electrical charge

$$\begin{cases} R_{BV} = R_{BV0}\, e^{\alpha_{BV}\tau} \\ \Omega_{BV} = \Omega_{BV0}\, e^{\beta_{BV}\tau} \end{cases} \qquad \begin{cases} q_{\lambda BV} = q_{\lambda BV0}\, e^{(2\alpha_{BV}+\beta_{BV})\tau} \\ M_0^{BV} = M_{00}^{BV}\, e^{-7\tau/8} \end{cases}$$

$$\begin{cases} \tau_{ext\ LD}^{BV}(r) \equiv -\dfrac{4G_{grav}}{c_t^2}\dfrac{M_0^{BV}}{r} \\ E_{dist}^{BV}(\tau^{externe}) \equiv M_0^{BV} c_t^2\, e^{\tau^{externe}/8} \\ M_0^{BV} \equiv \dfrac{1}{2c_t^2}(K_2+K_3)\zeta_{BV} R_{BV}^3 \Omega_{BV}^2 \equiv \left[\dfrac{1}{2c_{t0}^2}(K_2+K_3)\zeta_{BV} R_{BV0}^3 \Omega_{BV0}^2\right] e^{-7\tau_0/8} \end{cases}$$

The prismatic edge dislocation loop **(BC)**, source of flexion charge

$$\begin{cases} R_{BC} = R_{BC0}\, e^{\alpha_{BC}\tau} \\ \vec{B}_{BC} = \vec{B}_{BC0}\, e^{\beta_{BC}\tau} \\ q_{\theta BC} = q_{\theta BC0}\, e^{\beta_{BC}\tau} \end{cases}$$

$$q_{\theta BC} \begin{cases} >0 & \text{if vacancy loop} \\ <0 & \text{if interstitial loop} \end{cases}$$

$$M_{courbure}^{BC} \begin{cases} >0 & \text{if vacancy loop} \\ <0 & \text{if interstitial loop} \end{cases}$$

$$\begin{cases} \tau_{ext\ LD}^{BC}(r) \equiv -\dfrac{4G_{grav}}{c_t^2}\dfrac{M_{curvature}^{BC} + M_0^{BC}}{r} \\ E_{dist}^{BC}(\tau^{external}) \equiv M_0^{BC} c_t^2\, e^{(\alpha_{BC}+2\beta_{BC})\tau^{external}} \\ M_0^{BC} \equiv \left(\dfrac{K_2}{K_3}\right)^2 \dfrac{K_3\zeta_{BC} R_{BC} \vec{B}_{BC}^2}{c_t^2} \equiv \left[\left(\dfrac{K_2}{K_3}\right)^2 \dfrac{K_3\zeta_{BC} R_{BC0} \vec{B}_{BC0}^2}{c_{t0}^2}\right] e^{(\alpha_{BC}+2\beta_{BC}-1)\tau_0} \\ M_{curvature}^{BC} = -\dfrac{2\pi K_2}{5c_t^2} R_{BC}^2 \vec{n}\vec{B}_{BC} = \left[-\dfrac{2\pi K_2}{5c_{t0}^2} R_{BC0}^2 \vec{n}\vec{B}_{BC0}\right] e^{(2\alpha_{BC}+\beta_{BC}-1)\tau_0} \end{cases}$$

The slip mixed dislocation loop **(BM)**, source of an electric dipolar moment

$$\begin{cases} R_{BM} = R_{BM0}\, e^{\alpha_{BM}\tau} \\ \vec{B}_{BM} = \vec{B}_{BM0}\, e^{\beta_{BM}\tau} \end{cases}$$

$$\begin{cases} \tau_{ext\ LD}^{BM}(r) \equiv -\dfrac{4G_{grav}}{c_t^2}\dfrac{M_0^{BM}}{r} \\ E_{dist}^{BM}(\tau^{external}) \equiv M_0^{BM} c_t^2\, e^{(\alpha_{BM}+2\beta_{BM})\tau^{external}} \\ M_0^{BM} \equiv \dfrac{K_2+K_3}{4c_t^2}\zeta_{BM} R_{BM} \vec{B}_{BM}^2 \equiv \left[\dfrac{K_2+K_3}{4c_{t0}^2}\zeta_{BM} R_{BM0} \vec{B}_{BM0}^2\right] e^{(\alpha_{BM}+2\beta_{BM}-1)\tau_0} \end{cases}$$

The macroscopic vacancy **(L)**, analogous to a black hole

$$M_{grav}^{(L)}(\tau_0) \begin{cases} <0 & \Leftrightarrow & \tau_0 < -1 \\ >0 & \Leftrightarrow & -1 < \tau_0 < \tau_{0cr} \\ <0 & \Leftrightarrow & \tau_0 > \tau_{0cr} \end{cases}$$

$$\begin{cases} R_L = \sqrt[3]{\dfrac{3N_L}{4\pi n_0 e}} \\ \tau_{ext}^{(L)}(r) \equiv -\left(1+\tau_0+\tau^{externe}\right)\dfrac{R_L}{r} \equiv -\dfrac{4G_{grav}}{c_t^2}\dfrac{M_{grav}^{(L)}}{r} \\ E_{grav}^{(L)} \equiv M_{grav}^{(L)} c_t^2 \left(1+\dfrac{\tau^{externe}}{(1+\tau_0)}\right) \\ M_{grav}^{(L)} \equiv \dfrac{R_\infty^2}{c_t^2}\sqrt[3]{\dfrac{6\pi^2 N_L}{n_0 e}}\left(K_0 - 4K_2/3 - 2K_1(1+\tau_0)\right)(1+\tau_0) \\ \qquad \equiv \left[mR_{\infty 0}^2 \sqrt[3]{6\pi^2 e^{-1} n_0^2 N_L}\right] \dfrac{\left(K_0 - 4K_2/3 - 2K_1(1+\tau_0)\right)(1+\tau_0)}{K_2+K_3} e^{-\frac{\tau_0}{3}} \end{cases}$$

The macroscopic interstitial **(I)**, analogous to a neutron star

$$M_{grav}^{(I)}(\tau_0) \begin{cases} >0 & \Leftrightarrow & \tau_0 < \tau_{0cr} \\ <0 & \Leftrightarrow & \tau_0 > \tau_{0cr} \end{cases}$$

$$\begin{cases} R_I \equiv \sqrt[3]{\dfrac{3N_I}{4\pi n}} \equiv \sqrt[3]{\dfrac{3N_I}{4\pi n_0}} e^{\frac{\tau_0}{3}} \\ \tau_{ext}^{(I)}(r) \equiv -\dfrac{4G_{grav}}{c_t^2}\dfrac{M_{grav}^{(I)}}{r} \\ E_{grav}^{(I)} \equiv M_{grav}^{(I)} c_t^2\, e^{\tau^{external}} \\ M_{grav}^{(I)} \equiv \dfrac{N_I}{nc_t^2}\left(K_0 - 4K_2/3 - 2K_1(1+\tau_0)\right) \equiv mN_I \dfrac{K_0 - 4K_2/3 - 2K_1(1+\tau_0)}{K_2+K_3} \end{cases}$$

Figure 8.8 *- Dependence of the topological singularities of the lattice in the volume expansion of the lattice*

of the twist disclination loop, either $\alpha_{BC} = 1/4$, but it could also be similar to the dependence of the lattice pitch, either $\alpha_{BC} = 1/3$. As for the dependence of the Burgers vector, which must be a vector of the lattice, it should in principle take on value $\beta_{BC} = 1/3$. But for the moment, we will not make a choice of these values, and we will keep the parameters α_{BC} and β_{BC} , because the exact values of these parameters are not called to play a crucial role for the rest of the theory. The expressions of R_{BC} , of $\vec{\boldsymbol{B}}_{BC}$ and of $\boldsymbol{q}_{\theta BC}$ thus obtained make it possible to write four important relationships for a prismatic edge dislocation loop in the presence of a background expansion field τ_0 and / or an external expansion field $\tau^{external}$ generated by other singularities, remembering that the charge of curvature and the gravitational mass associated with the charge of curvature can be positive or negative according to the nature of the charge, as indicated in the table of figure 8.8.

In the case of a slip mixed dislocation loop ***(BM)***, we also do not know a priori the dependencies in volume expansion of the radius of the loop and its Burgers vector. We therefore introduce expansion constants α_{BM} and β_{BM} allowing us to write the relationships shown in figure 8.8.

We have seen, from the dilation of time and the curvature of the wave rays in a weak gravitational field, that the values which must take the parameters α_{BV} and β_{BV} associated with the twist disclination loops cannot be determined in a simple and unequivocal manner, but they must at least satisfy the relationship $3\alpha_{BV} + 2\beta_{BV} = 1/8$. For the edge and mixed dislocation loops, there are again no known experimental effects which could make it possible to simply deduce the value of the parameters $\alpha_{BC}, \beta_{BC}, \alpha_{BM}, \beta_{BM}$. The values of the different parameters $\alpha_{BV}, \beta_{BV}, \alpha_{BC}, \beta_{BC}, \alpha_{BM}, \beta_{BM}$ are therefore not known to us, except for the relationship $3\alpha_{BV} + 2\beta_{BV} = 1/8$. But the actual values of these parameters should most likely have physical consequences which should be accessible by the experience.

But there is yet another much more confusing physical consequence of these parameters. Indeed, if the radius of theedge loops were to depend on a value different from 1/3, therefore that $\alpha_{BC} \neq 1/3$, to ensure the existence of electron-like dispirations, this would mean that the number of lacunar or interstitial sites of the edge loop should vary if local expansion changes. However, the only possibility to vary this number of sites is that the edge loop behaves like a source or a pit of vacancies or interstitials in the presence of a variation of τ . And this effect would have very surprising consequences on Maxwell's equations, because in this case of divergent flow of vacancies or interstitials, Maxwell's equation $\mathrm{div}(n\vec{p}^{rot}) = 0$ should be replaced by $\mathrm{div}(n\vec{p}) \neq 0$, with for analog $\mathrm{div}\vec{B} \neq 0$ in electromagnetism. In other words, in the case where $\alpha_{BC} \neq 1/3$ *the edge dislocation loops behave like magnetic monopoles in the presence of variations in the expansion of the cosmological lattice*. However, *this effect should be observable and measurable by the* ***HS*** *observers*, who should measure a very weak monopole magnetic component linked to the particles containing edge loops, and caused by the local variation of the expansion *(which the* ***HS*** *cannot measure in principle).* And this monopole component of particles containing edge loops should necessarily exist *under the effect of the cosmological expansion of the universe*, which could open up an exciting chapter in our approach.

In the case of a macroscopic vacancy, one can take again the relations obtained in figure 7.7a and transform them so as to obtain the various relations reported in table 8.8 in the

$$F_{grav}^{BV-BV} \cong G_{grav} \frac{M_{0(1)}^{BV} M_{0(2)}^{BV}}{d^2}$$

$$F_{grav}^{BC-BC} \cong \left(\alpha_{BC} + 2\beta_{BC}\right) G_{grav} \frac{M_{curvature(1)}^{BC} M_{0(2)}^{BC} + M_{curvature(2)}^{BC} M_{0(1)}^{BC}}{d^2} + 2\left(\alpha_{BC} + 2\beta_{BC}\right) G_{grav} \frac{M_{0(1)}^{BC} M_{0(2)}^{BC}}{d^2}$$

$$F_{grav}^{BM-BM} \cong 2\left(\alpha_{BM} + 2\beta_{BM}\right) G_{grav} \frac{M_{0(1)}^{BM} M_{0(2)}^{BM}}{d^2}$$

$$F_{grav}^{BV-BC} \cong \frac{1}{2} G_{grav} \frac{M_{curvature}^{BC} M_0^{BV}}{d^2} + \left(\frac{1}{2} + 4\left(\alpha_{BC} + 2\beta_{BC}\right)\right) G_{grav} \frac{M_0^{BV} M_0^{BC}}{d^2}$$

$$F_{grav}^{BV-BM} \cong \left(\frac{1}{2} + 4\left(\alpha_{BM} + 2\beta_{BM}\right)\right) G_{grav} \frac{M_0^{BV} M_0^{BM}}{d^2}$$

$$F_{grav}^{BC-BM} \cong 4\left(\alpha_{BM} + 2\beta_{BM}\right) G_{grav} \frac{M_{curvature}^{BC} M_0^{BM}}{d^2} + 4\left(\alpha_{BC} + 2\beta_{BC} + \alpha_{BM} + 2\beta_{BM}\right) G_{grav} \frac{M_0^{BC} M_0^{BM}}{d^2}$$

$$F_{grav}^{BV-L} \cong \frac{1}{2} G_{grav} \frac{9+\tau_0}{1+\tau_0} \frac{M_0^{BV} M_{grav}^{(L)}}{d^2} \cong \frac{c_t^2}{8}\left(9+\tau_0\right) \frac{M_0^{BV} R_L}{d^2}$$

$$F_{grav}^{BC-L} \cong 4G_{grav} \frac{1}{1+\tau_0} \frac{M_{curvature}^{BC} M_{grav}^{(L)}}{d^2} + 4G_{grav} \frac{1+\left(\alpha_{BC} + 2\beta_{BC}\right)\left(1+\tau_0\right)}{1+\tau_0} \frac{M_0^{BC} M_{grav}^{(L)}}{d^2}$$

$$\cong c_t^2 \frac{M_{curvature}^{BC} R_L}{d^2} + c_t^2 \left[1 + \left(\alpha_{BC} + 2\beta_{BC}\right)\left(1+\tau_0\right)\right] \frac{M_0^{BC} R_L}{d^2}$$

$$F_{grav}^{BM-L} \cong 4G_{grav} \frac{1+\left(\alpha_{BM} + 2\beta_{BM}\right)\left(1+\tau_0\right)}{1+\tau_0} \frac{M_0^{BM} M_{grav}^{(L)}}{d^2} \cong c_t^2 \left[1 + \left(\alpha_{BM} + 2\beta_{BM}\right)\left(1+\tau_0\right)\right] \frac{M_0^{BM} R_L}{d^2}$$

$$F_{grav}^{BV-I} \cong \frac{9}{2} G_{grav} \frac{M_0^{BV} M_{grav}^{(I)}}{d^2} \cong \frac{3c_t^2}{4R_\infty^2} \frac{M_0^{BV} R_I^3}{d^2}$$

$$F_{grav}^{BC-I} \cong 4G_{grav} \frac{M_{curvature}^{BC} M_{grav}^{(I)}}{d^2} + 4G_{grav}\left(1 + \alpha_{BC} + 2\beta_{BC}\right) \frac{M_0^{BC} M_{grav}^{(I)}}{d^2}$$

$$\cong \frac{2c_t^2}{3R_\infty^2} \frac{M_{curvature}^{BC} R_I^3}{d^2} + \frac{2c_t^2}{3R_\infty^2}\left(1 + \alpha_{BC} + 2\beta_{BC}\right) \frac{M_0^{BC} R_I^3}{d^2}$$

$$F_{grav}^{BM-I} \cong 4G_{grav}\left(1 + \alpha_{BM} + 2\beta_{BM}\right) \frac{M_0^{BM} M_{grav}^{(I)}}{d^2} \cong \frac{2c_t^2}{3R_\infty^2}\left(1 + \alpha_{BM} + 2\beta_{BM}\right) \frac{M_0^{BM} R_I^3}{d^2}$$



$$F_{grav}^{L-L} \cong \frac{8G_{grav}}{\left(1+\tau_0\right)} \frac{M_{grav(1)}^{(L)} M_{grav(2)}^{(L)}}{d^2} \cong \frac{c_t^4\left(1+\tau_0\right)}{2G_{grav}} \frac{R_{L(1)} R_{L(2)}}{d^2}$$

$$F_{grav}^{I-I} \cong 2G_{grav} \frac{M_{grav(1)}^{(I)} M_{grav(2)}^{(I)}}{d^2} \cong \frac{c_t^4}{18 G_{grav} R_\infty^4} \frac{R_{I(1)}^3 R_{I(2)}^3}{d^2}$$

$$F_{grav}^{L-I} \cong 4G_{grav} \frac{2+\tau_0}{1+\tau_0} \frac{M_{grav}^{(L)} M_{grav}^{(I)}}{d^2} \cong \frac{c_t^4}{6R_\infty^2} \frac{2+\tau_0}{G_{grav}} \frac{R_L R_I^3}{d^2}$$

Figure 8.9 *- The set of possible gravitational interaction forces between the various topological singularities of the lattice*

presence of a background expansion field τ_0 and / or of an external expansion field $\tau^{external}$ generated by other singularities. We note that the gravitational mass of the vacancy has the property of changing sign for two values of the background volume expansion of the lattice, and we also remember that the macroscopic vacancy is the only singularity *which necessarily becomes a black hole* when the background volume expansion of the lattice satisfies the relationship $\tau_0 > 1 - \tau^{external}$.

In the case of a macroscopic interstitial, one can also take again the relations obtained in

figure 7.8 and transform them so as to obtain the various relations reported in table 8.8 in the presence of a background expansion field τ_0 and / or of an external expansion field $\tau^{external}$ generated by other singularities. We can see that the gravitational mass of the interstitial has

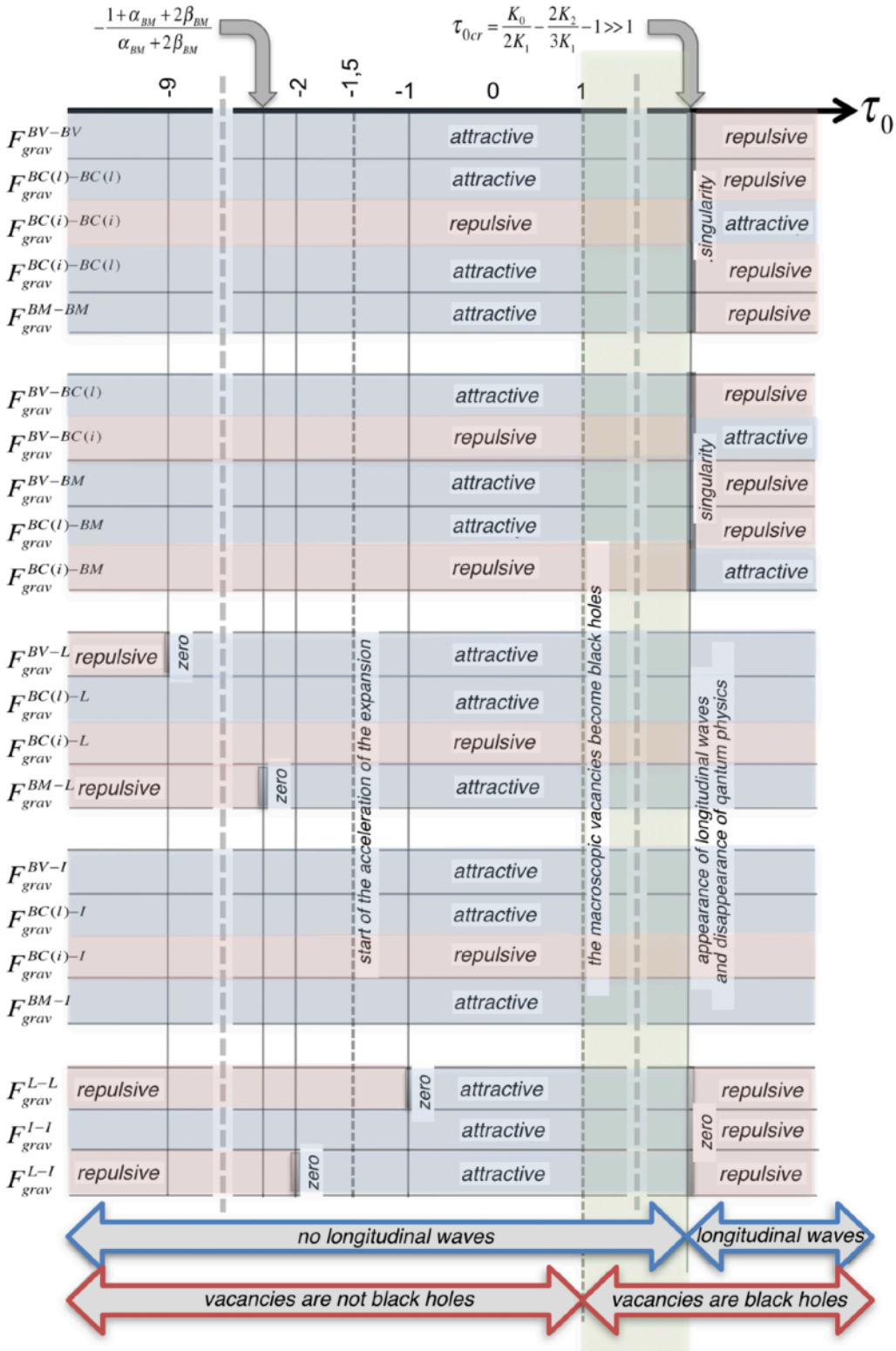


Figure 8.10 *- Attractive or repulsive behaviors of "gravitational" interaction forces between singularities as a function of the background expansion of the cosmological lattice*

the property of changing sign depending on the value of the background volume expansion of the lattice.

All of the gravitational interactions between the various singularities of the lattice

Using the relations obtained in table 8.8, we can then calculate the gravitational interaction forces that can appear between two of the singularities of the lattice, by calculating the energy increase of the two singularities by their interaction, as a function of the distance separating them. This increase therefore corresponds to the gravitational energy of interaction, and the derivative of this energy by the distance separating the two singularities corresponds to the force of gravitation between the two singularities. As five different topological singularities have been described in table 8.8, it is possible to calculate the gravitational interaction force in 15 different cases. By performing all these long and tedious calculations, we finally obtain all of these gravitational interaction forces as shown in figure 8.9. In this table, we see several important things:

- only the interaction between the twist disclination loops, carrying the "electric" charge of rotation, *exactly satisfies Newton's law of gravitation*. It should be noted that these loops have a much higher interaction energy than the other types of loops, and that *it is these loops which therefore very largely dominate the gravitational interactions between loops*,
- all the other interactions present a somewhat modified version of the gravitational interaction. Between the various loops, the interactions always depend on the "constant" of gravitation G_{grav}, but with an additional numerical factor which can contain the unknown parameters $\alpha_{BC}, \beta_{BC}, \alpha_{BM}, \beta_{BM}$,
- in the case where an edge dislocation loop occurs, there are always two interaction terms, one dependent on the mass of curvature $M_{curvature}^{BC}$ of the edge loop, and the other dependent on the mass of inertia M_0^{BC} of the edge loop. Since the curvature mass of the loop is much higher than its mass of inertia, the term containing the curvature mass largely dominates in the expression of the force of interaction. In addition, this dominant term can correspond to an attractive or negative interaction force since the mass of curvature of the corner loop is positive if the loop is of vacancy type and negative if the loop is of interstitial type,
- in the case where a macroscopic vacancy or a macroscopic interstitial intervenes, there are two possible formulations of the gravitational interaction force: the formulation which involves the masses of gravitation $M_{grav}^{(L)}$ and $M_{grav}^{(I)}$ and which resembles the formulation of Newton's law, but which has the disadvantage that the masses $M_{grav}^{(L)}$ and $M_{grav}^{(I)}$ strongly depend on the background expansion of the lattice, to the point of changing sign in certain domains of expansion. This is why we will prefer to use the second formulation which involves the radii R_L and R_I of the macroscopic singularities, and which has the advantage of being much simpler to analyze as regards the sign of the interaction (attractive or negative),
- on the basis of these expressions of "gravitational" interaction forces between singularities, we can deduce the attractive or repulsive behavior of all the interactions in table 8.8 as a function of the evolution of the background expansion τ_0 of the lattice. This behavior is reported in figure 8.10. In this figure, we did not respect the scale of values on the axis of expansion τ_0, especially with regard to the value of τ_{0cr} which is in fact extremely higher ($\tau_{0cr} >> 1$) since $K_0 >> K_1$ in the case of the perfect cosmological lattice.

We see that the gravitational interactions evolve strongly under the effect of the background expansion of the cosmological lattice. There appear in particular changes of sign of the interactions, which pass from the attractive mode to the repulsive mode or vice versa for certain values of expansion. These changes in sign of the interaction correspond either to *a zero crossing of the interaction force*, or to the appearance of *a singularity of the attraction force*.

It is obvious that figure 8.10 implies very important consequences on the cosmological behavior of the singularities of the lattice, that is to say on the evolution of the singularities during the cosmological expansion of the lattice. We will return to this subject in a future chapter. In figure 8.10, we have also reported certain phenomena associated with the evolution of the background expansion of the cosmological lattice, namely:

- first of all, for $\tau_0 = \tau_{0cr}$, the transition from the expansion domain without transverse waves, dominated by localized longitudinal eigen modes (to which we will return later), to the expansion domain where the real propagation of longitudinal waves appears, and where quantum physics disappears as we will see later,
- then the expansion domains in which the macroscopic vacancies are or are not black holes, with a transition for the expansion value $\tau_0 = 1$,
- the expansion value $\tau_0 = -3/2$ for which the evolution of the perfect cosmological lattice (figure 3.10) passes from the domain of the end of inflation (with decrease in the speed of expansion) to the domain of expansion with a speed of increasing expansion (acceleration of expansion),
- it is quite clear that the domain which is concerned by our current Universe lies between the values $\tau_0 = 1$ and $\tau_0 = \tau_{0cr}$, because it is in this domain that there are no longitudinal waves and that the macroscopic vacancies are black holes ,
- we will also notice that all the interactions which involve edge loops ***BC*** *(i)* of interstitial nature are repulsive in the domain between $\tau_0 = 1$ and $\tau_0 = \tau_{0cr}$, which will play a considerable role in the cosmological evolution of matter,
- finally, one did not represent in figure 8.10 the stages of the cosmological evolution which follow the critical value $\tau_0 = \tau_{0cr}$, namely the end of the acceleration of the expansion, the passage by a null value of the speed of expansion and finally the return to a contraction of the cosmological lattice (see figure 3.10).