Cosmological constant problems and their solutions

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Abstract

There are now two cosmological constant problems: (i) why the vacuum energy is so small and (ii) why it comes to dominate at about the epoch of galaxy formation. Anthropic selection appears to be the only approach that can naturally resolve both problems. The challenge presented by this approach is that it requires scalar fields with extremely flat potentials or four-form fields coupled to branes with an extremely small charge. Some recent suggestions are reviewed on how such features can arise in particle physics models.

I. THE PROBLEMS

Until recently, there was only one cosmological constant problem and hardly any solutions. Now, within the scope of a few years, we have made progress on both accounts. We now have two cosmological constant problems (CCPs) and a number of proposed solutions. In this talk I am going to review the situation, focussing mainly on the anthropic approach and on its implications for particle physics models. I realize that the anthropic approach has a low approval rating among physicists. But I think its bad reputation is largely undeserved. When properly used, this approach is quantitative and has no mystical overtones that are often attributed to it. Moreover, at present this appears to be the only approach that can solve both CCPs. I will also comment on other approaches to the problems.
The cosmological constant is (up to a factor) the vacuum energy density, $\rho_V$. Particle physics models suggest that the natural value for this constant is set by the Planck scale $M_P$,

$$\rho_V \sim M_P^4 \sim (10^{18} \text{ GeV})^4, \quad (1)$$

which is some 120 orders of magnitude greater than the observational bound,

$$\rho_V \lesssim (10^{-3} \text{ eV})^4. \quad (2)$$

In supersymmetric theories, one can expect a lower value,

$$\rho_V \sim \eta_{SUSY}^4, \quad (3)$$

where $\eta_{SUSY}$ is the supersymmetry breaking scale. However, with $\eta_{SUSY} \gtrsim 1 \text{ TeV}$, this is still 60 orders of magnitude too high. This discrepancy between the expected and observed values is the first cosmological constant problem. I will refer to it as the old CCP.

Until recently, it was almost universally believed that something so small could only be zero, due either to some symmetry or to a dynamical adjustment mechanism. (For a review of the early work on CCP, see [1].) It therefore came as a surprise when recent observations [2] provided evidence that the universe is accelerating, rather than decelerating, suggesting a non-zero cosmological constant. The observationally suggested value is

$$\rho_V \sim \rho_{M0} \sim (10^{-3} \text{ eV})^4, \quad (4)$$

where $\rho_{M0}$ is the present density of matter. This brings yet another puzzle. The matter density $\rho_M$ and the vacuum energy density $\rho_V$ scale very differently with the expansion of the universe, and there is only one epoch in the history of the universe when $\rho_M \sim \rho_V$. It is difficult to understand why we happen to live in this special epoch. Another, perhaps less anthropocentric statement of the problem is why the epoch when the vacuum energy starts dominating the universe ($z_V \sim 1$) nearly coincides with the epoch of galaxy formation ($z_G \sim 1 - 3$), when the giant galaxies were assembled and the bulk of star formation has occurred:
\( t_V \sim t_G. \)  

This is the so-called cosmic coincidence problem, or the second CCP.

II. PROPOSED SOLUTIONS

A. Quintessence

Much of the recent work on CCP involves the idea of quintessence \(^3\)). Quintessence models require a scalar field \( \phi \) with a potential \( V(\phi) \) approaching zero at large values of \( \phi \). A popular example is an inverse power law potential,

\[ V(\phi) = M^{4+\beta} \phi^{-\beta}, \]

with a constant \( M \ll M_P \). It is assumed that initially \( \phi \ll M_P \). Then it can be shown that the quintessence field \( \phi \) approaches an attractor “tracking” solution

\[ \phi(t) \propto t^{2/(2+\beta)} \]

in which its energy density grows relative to that of matter,

\[ \rho_\phi / \rho_M \sim \phi^2 / M_P^2. \]

When \( \phi \) becomes comparable to \( M_P \), its energy dominates the universe. At this point the nature of the solution changes: the evolution of \( \phi \) slows down and the universe enters an epoch of accelerated expansion. The mass parameter \( M \) can be adjusted so that this happens at the present epoch.

A nice feature of the quintessence models is that their evolution is not sensitive to the choice of the initial conditions. However, I do not think that these models solve either of the two CCPs. The potential \( V(\phi) \) is assumed to vanish in the asymptotic range \( \phi \rightarrow \infty \). This assumes that the old CCP has been solved by some unspecified mechanism. The coincidence problem also remains unresolved, because the time of quintessence domination depends on the choice of the parameter \( M \), and there seems to be no reason why this time should coincide with the epoch of galaxy formation.
B. k-essence

A related class of models involves k-essence, a scalar field with a non-trivial kinetic term in the Lagrangian [4],

\[ L = \phi^{-2}K[(\nabla \phi)^2]. \]  

(9)

For a class of functions \( K(X) \), the energy density of k-essence stays at a constant fraction of the radiation energy density during the radiation era,

\[ \rho_\phi / \rho_{rad} \approx \text{const}, \]  

(10)

and starts acting as an effective cosmological constant with the onset of matter domination. The function \( K(X) \) can be designed so that the constant in Eq. (10) is \( \lesssim 10^{-2} \), thus avoiding conflict with nucleosynthesis, and that k-essence comes to dominate at \( z \sim 1 \).

This is an improvement over quintessence, since the accelerated expansion in this kind of models always begins during the matter era. Galaxy formation can also occur only in the matter era, but still there seems to be no reason why the two epochs should coincide. The epoch of k-essence domination \( z_V \) is determined by the form of the function \( K(X) \), and the epoch of galaxy formation \( z_G \) is determined by the amplitude of primordial density fluctuations,

\[ Q = \delta \rho / \rho \sim 10^{-5}. \]  

(11)

It is not clear why these seemingly unrelated quantities should give \( z_V \sim z_G \) within one order of magnitude. And of course the old CCP also remains unresolved.

C. A small cosmological constant from fundamental physics

One possibility here is that some symmetry of the fundamental physics requires that the cosmological constant should be zero. A small value of \( \rho_V \) could then arise due to a small violation of this symmetry. One could hope that \( \rho_V \) would be given by an expression like
\[ \rho_V \sim M_W^8/M_P^4 \sim (10^{-3} \text{ eV})^4, \]  

(12)

where \( M_W \sim 10^3 \) GeV is the electroweak scale. There have been attempts in this direction \footnote{5}, but no satisfactory implementation of this program has yet been developed. And even if we had one, the time coincidence \( t_V \sim t_G \) would still remain a mystery.

Essentially the same remarks apply to the braneworld \footnote{6} and holographic \footnote{7} approaches to CCPs.

D. Anthropic approach

According to this approach, what we perceive as the cosmological constant is in fact a stochastic variable which varies on a very large scale, greater than the present horizon, and takes different values in different parts of the universe. We shall see that situation of this sort can naturally arise in the context of the inflationary scenario.

The key observation here is that the gravitational clustering that leads to galaxy formation effectively stops at \( z \sim z_V \). An anthropic bound on \( \rho_V \) can be obtained by requiring that it does not dominate before the redshift \( z_{\text{max}} \) when the earliest galaxies are formed. With \( z_{\text{max}} \sim 5 \) one obtains \footnote{8}

\[ \rho_V \lesssim 200 \rho_M. \]  

(13)

For negative values of \( \rho_V \), a lower bound can be obtained by requiring that the universe does not recollapse before life had a chance to develop \footnote{9},

\[ \rho_V \gtrsim -\rho_M. \]  

(14)

The bound (13) is a dramatic improvement over (1) or (3), but it still falls short of the observational bound by a factor of about 50. If all values in the anthropic range (13) were equally probable, then \( \rho_V \sim \rho_M \) would still be ruled out at a 95% confidence level. However, the values in this range are not equally probable. The anthropic bound (13) specifies the value of \( \rho_V \) which makes galaxy formation barely possible. Most of the galaxies will be not
in regions characterized by these marginal values, but rather in regions where $\rho_V$ dominates after the bulk of galaxy formation has occurred, that is $z_V \lesssim 1$. 

This can be made quantitative by introducing the probability distribution as

$$d\mathcal{P}(\rho_V) = \mathcal{P}_*(\rho_V) \nu(\rho_V) d\rho_V.$$  

Here, $\mathcal{P}_*(\rho_V) d\rho_V$ is the prior distribution, which is proportional to the volume of those parts of the universe where $\rho_V$ takes values in the interval $d\rho_V$, and $\nu(\rho_V)$ is the average number of galaxies that form per unit volume with a given value of $\rho_V$. The calculation of $\nu(\rho_V)$ is a standard astrophysical problem; it can be done, for example, using the Press-Schechter formalism.

The distribution (15) gives the probability that a randomly selected galaxy is located in a region where the effective cosmological constant is in the interval $d\rho_V$. If we are typical observers in a typical galaxy, then we should expect to observe a value of $\rho_V$ somewhere near the peak of this distribution.

The prior distribution $\mathcal{P}_*(\rho_V)$ should be determined from the inflationary model of the early universe. Weinberg has argued that a flat distribution, $\mathcal{P}_*(\rho_V) = \text{const}$, (16) should generally be a good approximation. The reason is that the function $\mathcal{P}_*(\rho_V)$ is expected to vary on some large particle physics scale, while we are only interested in its values in the tiny anthropically allowed range. Analysis shows that this Weinberg’s conjecture is indeed true in a wide class of models, but one finds that it is not as automatic as one might expect.

Martel, Shapiro and Weinberg (see also [11,13]) presented a detailed calculation of $d\mathcal{P}(\rho_V)$ assuming a flat prior distribution. They found that the peak of the resulting probability distribution is close to the observationally suggested values of $\rho_V$.

The cosmic time coincidence is easy to understand in this approach. Regions of the universe where $t_V \ll t_G$ do not form any galaxies at all, whereas regions where $t_V \gg t_G$
are suppressed by “phase space”, since they correspond to a very tiny range of $\rho_V$. It was shown in Ref. [18] that the probability distribution for $t_G/t_V$ is peaked at $t_G/t_V \approx 1.5$, and thus most observers will find themselves in galaxies formed at $t_G \sim t_V$.

We thus see that the anthropic approach naturally resolves both CCPs. All one needs is a particle physics model that would allow $\rho_V$ to take different values and an inflationary cosmological model that would give a more or less flat prior distribution $P_*(\rho_V)$ in the anthropic range (13).

III. MODELS WITH A VARIABLE $\rho_V$

A. Scalar field with a very flat potential

One possibility is that what we perceive as a cosmological constant is in fact a potential $V(\phi)$ of some field $\phi(x)$ [13]. The slope of the potential is assumed to be so small that the evolution of $\phi$ is slow on the cosmological time scale. This is achieved if the slow roll conditions

\begin{align}
M^2_P V'' &\ll V \lesssim \rho_{M_0}, \\
M_P V' &\ll V \lesssim \rho_{M_0},
\end{align}

are satisfied up to the present time. These conditions ensure that the field is overdamped by the Hubble expansion, and that the kinetic energy is negligible compared with the potential energy (so that the equation of state is basically that of a cosmological constant term.) The field $\phi$ is also assumed to have negligible couplings to all fields other than gravity.

Let us now suppose that there was a period of inflation in the early universe, driven by the potential of some other field. The dynamics of the field $\phi$ during inflation are strongly influenced by quantum fluctuations, causing different regions of the universe to thermalize with different values of $\phi$. Spatial variation of $\phi$ is thus a natural outcome of inflation.
The probability distribution $\mathcal{P}_*(\phi)$ is determined mainly by the interplay of two effects. The first is the “diffusion” in the field space caused by quantum fluctuations. The dispersion of $\phi$ over a time interval $\Delta t$ is

$$\Delta \phi \sim H(H \Delta t)^{1/2}, \quad (19)$$

where $H$ is the inflationary expansion rate. The effect of diffusion is to make all values of $\phi$ equally probable over the interval $\Delta \phi$. The second effect is the differential expansion. Although $V(\phi)$ represents only a tiny addition to the inflaton potential, regions with larger values of $V(\phi)$ expand slightly faster, and thus the probability for higher values of $V(\phi)$ is enhanced. The time it takes the field $\phi$ to fluctuate across the anthropic range $\Delta \phi_{\text{anth}} \sim \rho_{M_0}/V'$ is $\Delta t_{\text{anth}} \sim (\Delta \phi_{\text{anth}})^2/H^3$, and the characteristic time for differential expansion is $\Delta t_{\text{de}} \sim H M_p^2/V$.

The effect of differential expansion is negligible if $\Delta t_{\text{anth}} \ll \Delta t_{\text{de}}$. The corresponding condition on $V(\phi)$ is

$$V'^2 \gg \rho_{M_0}^3/H^3 M_p^2. \quad (20)$$

In this case, the probability distribution for $\phi$ is flat in the anthropic range,

$$\mathcal{P}_*(\phi) = \text{const}. \quad (21)$$

The probability distribution for the effective cosmological constant $\rho_V = V(\phi)$ is given by

$$\mathcal{P}_*(\rho_V) = \frac{1}{V'} \mathcal{P}_*(\phi),$$

and it will also be very flat, since $V'$ is typically almost constant in the anthropic range. As we discussed in Section II, a flat prior distribution for the effective cosmological constant in the anthropic range entails an automatic explanation for the two cosmological constant puzzles.

On the other hand, if the condition (20) is not satisfied, then the prior probability for the field values with a higher $V(\phi)$ would be exponentially enhanced with respect to the
field values at the lower anthropic end. This would result in a prediction for the effective cosmological constant which would be too high compared with observations.

A simple example is given by a potential of the form

\[ V(\phi) = \rho_\Lambda + \frac{1}{2} \mu^2 \phi^2, \tag{22} \]

where \( \rho_\Lambda \) represents the "true" cosmological constant. We shall assume that \( \rho_\Lambda < 0 \), so that the two terms in (22) partially cancel one another in some parts of the universe. With \( |\rho_\Lambda| \sim (1 \text{ TeV})^4 \), the slow roll conditions (17), (18) give

\[ \mu < 10^{-90} M_P. \tag{23} \]

Thus, an exceedingly small mass scale must be introduced.

The condition (24) yields a lower bound on \( \mu \),

\[ \mu \gtrsim 10^{-137} M_P. \tag{24} \]

Here, I have used the upper bound on the expansion rate at late stages of inflation, \( H \lesssim 10^{-5} M_P \), which follows from the CMB observations.

We thus see that models with a variable \( \rho_V \) can be easily constructed in the framework of inflationary cosmology. The challenge here is to explain the very small mass scale (23) in a natural way.

**B. Four-form models**

Another class of models, first discussed by Brown and Teitelboim [20], assumes that the cosmological constant is due to a four-form field [21],

\[ F^{\alpha\beta\gamma\delta} = F^{e^{\alpha\beta\gamma\delta}}. \tag{25} \]

The field equation for \( F \) is \( \partial_\mu F = 0 \), so \( F \) is a constant, but it can change its value through nucleation of bubbles bounded by domain walls, or branes. The total vacuum energy density is given by
\[ \rho_V = \rho_\Lambda + F^2/2 \]  

(26)  

and once again it is assumed that \( \rho_\Lambda < 0 \). The change of the field across the brane is  

\[ \Delta F = q, \]  

(27)  

where the “charge” \( q \) is a constant fixed by the model. Thus, \( F \) takes a discrete set of values, and the resulting spectrum of \( \rho_V \) is also discrete. The four-form model has recently attracted much attention \([22-25,16]\) because four-form fields coupled to branes naturally arise in the context of string theory.

In the range where the bare cosmological constant is almost neutralized, \( |F| \approx |2\rho_\Lambda|^{1/2} \), the spectrum of \( \rho_V \) is nearly equidistant, with a separation  

\[ \Delta \rho_V \approx |2\rho_\Lambda|^{1/2}q. \]  

(28)  

In order for the anthropic explanation to work, \( \Delta \rho_V \) should not exceed the present matter density,  

\[ \Delta \rho_V \lesssim \rho_{m0} \sim (10^{-3} \text{ eV})^4. \]  

(29)  

With \( \rho_\Lambda \gtrsim (1 \text{ TeV})^4 \), it follows that  

\[ q \lesssim 10^{-90} M_P^2. \]  

(30)  

Once again, the challenge is to find a natural explanation for such very small values of \( q \).

In order to solve the cosmological constant problems, we have to require in addition that  

(i) the probability distribution for \( \rho_V \) at the end of inflation is nearly flat, \( P_*(\rho_V) \approx \text{const} \),  

and (ii) the brane nucleation rate is sufficiently low, so that the present vacuum energy does not drop significantly in less than a Hubble time. Models satisfying all the requirements can be constructed, but the conditions (i), (ii) significantly constrain the model parameters. For a detailed discussion, see [16].
IV. EXPLAINING THE SMALL PARAMETERS

Both scalar field and four-form models discussed above have some seemingly unnatural features. The scalar field models require extremely flat potentials and the four-form models require branes with an exceedingly small charge. The models cannot be regarded as satisfactory until the smallness of these parameters is explained in a natural way. Here I shall briefly review some possibilities that have been suggested in the literature.

A. Scalar field renormalization

Let us start with the scalar field model. Weinberg [14] suggested that the flatness of the potential could be due to a large field renormalization. Consider the Lagrangian of the form

\[ L = \frac{Z}{2} (\nabla \phi)^2 - V(\phi). \] (31)

The potential for the canonically normalized field \( \phi' = \sqrt{Z} \phi \) will be very flat if the field renormalization constant is very large, \( Z \gg 1 \).

More generally, the effective Lagrangian for \( \phi \) will include non-minimal kinetic terms [23, 16],

\[ L = \frac{1}{2} F^2(\phi)(\nabla \phi)^2 - V(\phi). \] (32)

Take for example \( F = e^{\phi/M} \). Then the potential for the canonical field \( \psi = M e^{\phi/M} \) is \( V(M \ln(\psi/M)) \). This will typically have a very small slope if \( V(\phi) \) is a polynomial function. It would be good to have some particle physics motivation either for a large running of the field renormalization, or for an exponential function \( F(\phi) \) in the Lagrangian [32].

B. A discrete symmetry

Another approach attributes the flatness of the potential to a spontaneously broken discrete symmetry [28]. The main ingredients of the model are: (1) a four-form field \( F_{\mu\nu\sigma\tau} \)
which can be obtained from a three-form potential, \( F_{\mu \nu \sigma \tau} = \partial_{[\mu} A_{\nu \sigma \tau]} \), (2) a complex field \( X \) which develops a vacuum expectation value

\[
\langle X \rangle = \eta e^{ia}
\]  

(33)

and whose phase \( a \) becomes a Goldstone boson, and (3) a scalar field \( \Phi \) which is used to break a discrete \( Z_{2N} \) symmetry.

The action is assumed to be invariant under the following three symmetries: (1) \( Z_{2N} \) symmetry under which

\[
\Phi \to \Phi e^{i\pi/N}, \quad a \to -a \quad (\text{or} \quad X \to X^\dagger),
\]

(34)

(2) a symmetry of global phase transformations

\[
a \to a + \text{const},
\]

(35)

and (3) the three-form gauge transformation

\[
A_{\mu \nu \sigma} \to A_{\mu \nu \sigma} + \partial_{[\mu} B_{\nu \sigma]},
\]

(36)

where \( B_{\nu \sigma} \) is a two-form. Below the symmetry breaking scales of \( X \) and \( \Phi \), the effective Lagrangian for the model can be written as

\[
L = \eta^2 (\partial_{\mu} a)^2 - \frac{1}{4} F^2 + (\text{effective interactions}).
\]

(37)

The interactions generally include all possible terms that are compatible with the symmetries. Among such terms is the mixing of the Goldstone \( a \) with the three-form potential,

\[
g \eta^2 \frac{(\Phi)^N}{M_P^2} \epsilon^{\nu \sigma \tau} A_{\nu \sigma \tau} \partial_{\mu} a,
\]

(38)

where \( g \lesssim 1 \) is a dimensionless coupling and I have assumed that the Planck scale \( M_P \) plays the role of the ultraviolet cutoff of the theory.

The effect of the mixing term (38) is to give a mass

\[
\mu = g \eta \frac{(\Phi)^N}{M_P^2}
\]

(39)
to the field $a$. This mass can be made very small if $\langle \Phi \rangle \ll M_P$ and $N$ is sufficiently large. For example, with $\langle \Phi \rangle \sim 1 \text{ TeV}$, $\eta \ll M_P$ and $N \geq 6$, we have $\mu \ll 10^{-90} M_P$, as required.

Models of this type can also be used to generate branes with a very small charge. In this case $a$ is assumed to be a pseudo-Goldstone boson, like the axion, and the theory has domain wall solutions with $a$ changing by $2\pi$ across the wall. The mixing of $a$ and $A$ couples these walls to the four-form field, and it can be shown that the corresponding charge is

$$q = 2\pi g \eta^2 \frac{\langle \Phi \rangle^N}{M_P^N}. \quad (40)$$

Once again, the anthropic constraint on $q$ is satisfied for $\langle \Phi \rangle \sim 1 \text{ TeV}$, $\eta \ll M_P$ and $N \geq 6$.

The central feature of this approach is the $Z_{2N}$ symmetry (34). What makes this symmetry unusual is that the phase transformation of $\Phi$ is accompanied by a charge conjugation of $X$. It can be shown, however, that such a symmetry can be naturally embedded into a left-right symmetric extension of the standard model [26].

C. String theory inspired ideas

Feng et. al. [24] have argued that branes with extremely small charge and tension can naturally arise due to non-perturbative effects in string theory. A potential problem with this approach is that the small brane tension and charge appear to be unprotected against quantum corrections below the supersymmetry breaking scale [26]. The cosmology of this model is also problematic, since it is hard to stabilize the present vacuum against copious brane nucleation [16].

A completely different approach was taken by Bousso and Polchinski [22]. They assume that several four-form fields $F_i$ are present so that the vacuum energy is given by

$$\rho_V = \rho_\Lambda + \frac{1}{2} \sum_i F_i^2. \quad (41)$$

The corresponding charges $q_i$ are not assumed to be very small, but Bousso and Polchinski have shown that with multiple four-forms the spectrum of the allowed values of $\rho_V$ can be
sufficiently dense to satisfy the anthropic condition \((29)\) in the range of interest. However, the situation here is quite different from that in the single-field models. The vacua with nearby values of \(\rho_V\) have very different values of \(F_i\), and there is no reason to expect the prior probabilities for these vacua to be similar. Moreover, the low energy physics in different vacua is likely to be different, so the process of galaxy formation and the types of life that can evolve will also differ. It appears therefore that the anthropic approach to solving the cosmological constant problems cannot be applied to this case \([25]\).

V. CONCLUDING REMARKS

In conclusion, it appears that the only approach that can solve both cosmological constant problems is the one that attributes them to anthropic selection effects. In this approach what we perceive as the cosmological constant is in fact a stochastic variable which varies from one part of the universe to another. A typical observer then finds himself in a region with a small cosmological constant which comes to dominate at about the epoch of galaxy formation. Cosmological models of this sort can easily be constructed in the framework of inflation. What one needs is either a scalar field with a very flat potential, or a four-form field coupled to branes with a very small charge. Some interesting suggestions have been made on how such features can arise; the challenge here is to implement these suggestions in well motivated particle physics models. (One attempt in this direction has been made in \([26]\).)

There are also problems to be addressed on the astrophysical side of the anthropic approach. All anthropic calculations of the probability distribution \((15)\) for \(\rho_V\) assumed that observers are in giant galaxies like ours and identified \(\nu(\rho_V)\) in Eq. \((17)\) with the density of such galaxies. This, however, needs some justification\(^{1}\). In the hierarchical structure formation scenario, dwarf galaxies could form as early as \(z = 10\), and if they are included among

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\(^{1}\)I am grateful to David Spergel for emphasizing this to me.
the possible sites for observers, then the expected epoch of vacuum domination would be $z_V \sim 10$.

One problem with dwarf galaxies is that if the mass of a galaxy is too small, then it cannot retain the heavy elements dispersed in supernova explosions. Numerical simulations suggest that the fraction of heavy elements retained is $\sim 30\%$ for a $10^9 M_\odot$ galaxy and is negligible for much smaller galaxies [27]. The heavy elements are necessary for the formation of planets and of observers, and thus one has to require that the structure formation hierarchy should evolve up to mass scales $\sim 10^9 M_\odot$ or higher prior to vacuum domination.

Another point to note is that smaller galaxies formed at earlier times have a higher density of matter. If this translates into a higher density of stars (or dark matter clumps), then we may have additional constraints by requiring that the timescales for disruption of planetary orbits by stellar encounters should not be too short. However, the cross-section for planetary orbit disruption is rather small (comparable to the size of the Solar system), and since close stellar encounters are quite rare in our galaxy, one does not expect a large effect from a modest density enhancement in dwarf galaxies.

An interesting possibility is that disruption of orbits of comets, rather than planets, could be the controlling factor of anthropic selection [28]. Comets move around the Sun, forming the Oort cloud of radius $\sim 0.1$ pc (much greater than the Solar system!). Whenever a star or a molecular cloud passes by, the orbits of some comets are disrupted and some of them enter the interior of the Solar system. Occasionally they hit planets, causing mass extinctions. The time it took to evolve intelligent beings after the last major hit is comparable to the typical time interval between hits on Earth ($\sim 10^8$ yrs), so one could argue that a substantial increase in the rate of hits might interfere with the evolution of observers. There are, of course, quite a few blanks to be filled in this scenario, and at present we are far from being able to reliably quantify the scale of bound systems to be used in the definition of $\nu(\rho_V)$. However, if the anthropic approach is on the right track, then one can predict that future research will show the relevant scale to be that of giant galaxies.

Finally, I would like to mention the possibility of a ‘compromise’ solution to CCPs.
It is conceivable that the cosmological constant will eventually be determined from the fundamental theory. For example, it could be given by the relation (12). This would solve the old CCP. The time coincidence problem could then be solved anthropically if the amplitude of density fluctuations $Q$ is a stochastic variable. With some mild assumptions about the probability distribution $P_s(Q)$, one finds that most galaxies will form at about the time of vacuum domination [10].

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