Primordial black holes from fifth forces

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Primordial black holes can be produced by a long range attractive fifth force stronger than gravity, mediated by a light scalar field interacting with non-relativistic "heavy" particles. As soon as the energy fraction of heavy particles reaches a threshold, the fluctuations rapidly grow non-linear. The overdensities collapse into black holes or similar screened objects, without the need of any particular feature in the spectrum of primordial density fluctuations generated during inflation. We discuss if such primordial black holes can constitute the total dark matter component in the Universe.

arXiv:1711.09915v1 [astro-ph.CO] 27 Nov 2017

In spite of the many observations leading to the establishment of dark matter as an essential ingredient of modern cosmology, its fundamental nature remains an open question. Among the many dark matter candidates, primordial black holes (BH) [1–4] are interesting since they could account for the gravitational wave signals observed by LIGO and VIRGO [5] (see also [6] and references therein). Several production mechanisms for primordial BH have been proposed [7, 8].

In this paper we present a novel framework for primordial BH formation which does not rely on a particular feature in the spectrum of primordial density fluctuations generated during inflation. The main assumption of our scenario is the presence in the early universe of a long-range interaction stronger than gravity. We associate this *fifth force* to a light scalar field interacting with some heavy degrees of freedom beyond the Standard Model particle content. More precisely, we assume that during some epoch in cosmology the Hubble parameter H is larger than the mass of a scalar field ϕ. If this scalar field couples to some "heavy particles" ψ with masses larger than H, it mediates an attractive fifth force which is effectively long range, similar to gravity. This attraction can be, however, substantially stronger than the gravitational attraction. As a result, the fluctuations in the energy density of the heavy fields can grow rapidly and eventually become non-linear. If the range and strength of the fifth force is large enough, it seems likely that a substantial part of the ψ fluid will collapse into BHs or similar screened objects.

This BH formation process could occur very early in cosmology, for example nearly after the end of inflation, with the ϕ and ψ fields associated to models within a grand unified framework or similar. For different properties of the participating particles it could also take place rather late in the cosmological history, say after nucleosynthesis. The heavy particles remaining outside primordial BHs might decay after the formation epoch and be unobservable today. The scalar field could relax after BH formation to a minimum of its effective potential with mass eventually exceeding the decreasing Hubble parameter. In this case the field ϕ would not be observable at present either. Alternatively, ϕ could be an additional dark matter candidate, or have a runaway behavior and be associated with dynamical dark energy. The BH formation process is not affected by what happens to the participating fields or particles at later times. Once BHs are formed, they behave as non-relativistic matter. If the total energy density of BHs is large enough, they could constitute the dark matter component of our Universe.

Fifth-force interactions. Consider the action

$$S=\int d^4x\sqrt{-g}\left[\frac{M_P^2}{2}R+\mathcal{L}_R+\mathcal{L}(\phi)+\mathcal{L}(\phi,\psi)\right]\,,\quad (1)$$

with $M_P=(8\pi G)^{-1/2}=2.435\times10^{18}\,\mathrm{GeV}$ the reduced Planck mass and $\mathcal{L}_R$ a radiation component that we assume to dominate the background evolution during the epoch relevant for primordial BH formation. The term

$$\mathcal{L}(\phi)=-\frac{1}{2}\partial^\mu\phi\partial_\mu\phi-V(\phi)\,,\quad (2)$$

stands for the Lagrangian density of an almost massless scalar field ϕ. We will assume for simplicity that the potential $V(\phi)$ can be neglected during the epoch of BH formation such that ϕ becomes effectively a massless field. This approximation is justified if the scalar field mass is smaller than H, both for a potential with a minimum or for a runaway potential. For definiteness we take the heavy particle ψ to be a fermion. The interactions with the field ϕ arise via a field-dependent mass term $m(\phi)$,

$$\mathcal{L}(\phi,\psi)=i\bar{\psi}\left(\gamma^\mu\nabla_\mu-m(\phi)\right)\psi\,,\quad (3)$$

for example by a Yukawa type coupling $\sim g\phi\bar{\psi}\psi$.

We will assume that the field equations derived from the action (1) admit a perfect fluid description and consider a flat Friedmann-Lemaître-Robertson-Walker Universe. The background evolution equations for the average ϕ and ψ energy densities read

$$\dot\rho_\phi+3H(\rho_\phi+p_\phi)=\frac{\beta}{M_P}\left(\rho_\psi-3p_\psi\right)\dot\phi\,,\quad (4)$$

$$\dot\rho_\psi+3H\left(\rho_\psi+p_\psi\right)=-\frac{\beta}{M_P}\left(\rho_\psi-3p_\psi\right)\dot\phi\,,\quad (5)$$

with $\rho_\phi=p_\phi=\dot\phi^2/2$, $H^2=\rho/(3M_P^2)$ and $\rho=\rho_R+\rho_\phi+\rho_\psi$. The ϕ and ψ fluids are coupled whenever the

ψ particles are non-relativistic ($\rho_\psi \neq 3p_\psi$, with p_ψ the ψ-fluid pressure). The coupling function $\beta(\phi)$ measures the dependence of the effective mass $m(\phi)$ on the scalar field ϕ,

$$\beta(\phi) = -M_P \frac{\partial \ln m(\phi)}{\partial \phi} \,. \tag{6}$$

Its normalization involving M_P has been chosen such that for $\beta^2 = 1/2$ the scalar-field mediated attraction has the same strength as gravity. The combined strength of the fifth force and gravity is proporportional to

$$Y \equiv 1 + 2\beta^2 \,. \tag{7}$$

This type of scenarios has been extensively studied in the literature [9–14], but not within the context of primordial black formation. A given model is specified by a choice of $\beta(\phi)$. The value of β can be rather large. Consider for instance a renormalizable interaction term of the form $m(\phi)\bar{\psi}\psi = m_0\bar{\psi}\psi + g\phi\bar{\psi}\psi$ with m_0 a constant mass parameter and g a dimensionless coupling. For this particular example, we can rewrite Eq. (6) as $\beta(\phi) = -g\, M_P/m(\phi)$, which leads to $|\beta| \gg 1$ if $m(\phi)/M_P \ll g$. Even a small value of the Yukawa coupling g can be largely overwhelmed by the ratio M_P/m. The factor $2\beta^2$ can be therefore naturally rather large. For illustration purposes we will neglect the field dependence of β in the following sections. The qualitative features of the BH formation scenario presented here do not rely on this approximation. They also hold if we take for ψ a non-relativistic bosonic particle rather than a fermion.

Background cosmology. For a non-relativistic fluid of ψ particles, $p_\psi = 0$, the set of equations (4)-(5) admits a scaling solution. In the limit $\beta \gg 1$ it reads [15]

$$\Omega_\psi = \frac{1}{3\beta^2}\,, \qquad \Omega_\phi = \frac{1}{6\beta^2}\,, \qquad \Omega_R = 1 - \frac{1}{2\beta^2}\,, \tag{8}$$

where $\Omega_i \equiv \rho_i/(3M_P^2H^2)$ and $i = R, \phi, \psi$. During this scaling phase, the fermion energy density *tracks* the background component, $\rho_\psi \sim \rho_R \sim a^{-4}$. Combining this scaling with the intuitive solution of Eq. (5), $\rho_\psi = \rho_{\psi,0}\, a^{-3} \exp\left(-\beta\phi/M_P\right)$, we obtain a relation between β and the variation of the ϕ field during the scaling phase, namely

$$\phi' = M_P/\beta \,, \tag{9}$$

with primes denoting derivatives with respect to the number of e-folds $dN \equiv d\ln a$, e.g. $\phi' = \dot{\phi}/H$.

Alternatively, Ω_ψ could be even smaller than $1/(3\beta^2)$. The fifth force plays then no role for the background evolution and ρ_ψ decays as non-relativistic matter, $\rho_\psi \sim a^{-3}$. In consequence, the density parameter Ω_ψ increases with time until it reaches the scaling solution $\Omega_\psi = 1/(3\beta^2)$ at some time $t_{\rm in}$. The evolution of the different energy densities is depicted in Fig. 1.

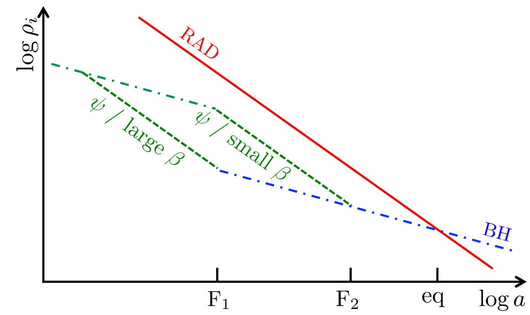

Figure 1. Schematic evolution of the energy density of different species. The red and blue lines stand respectively for radiation and BH dark matter. Green lines denotes ρ_ψ. Larger values of β imply earlier BH formation and therefore smaller masses. For $\Omega_{\rm BH}(a_{\rm eq}) < 1/2$ the radiation line moves upwards.

Which of the above regimes is realized during the BH formation period depends on the initial conditions for the energy densities ρ_ψ and ρ_ϕ, which are typically set at the end of inflation. For the large β^2 values we will be interested in, both the difference between various initial conditions and the question if the scaling solution (8) is reached or not become unimportant from the point of view of the background evolution. For all practical purposes, the background behaves as a standard RD Universe with a Hubble parameter obeying $\mathcal{H}' = -\mathcal{H}$ with $\mathcal{H} = aH$.

Growth of fluctuations. Due to the strong attractive force for $Y \gg 1$ and the decreasing particle masses, the ψ fluctuations grow rapidly for sufficiently large Ω_ψ. On scales k^{-1} sufficiently inside the horizon, the linear perturbations δ_ψ in the ψ fluid evolve as [16]

$$\delta_\psi'' + \left(1 + \frac{\mathcal{H}'}{\mathcal{H}} - \frac{\beta\phi'}{M_P}\right)\delta_\psi' - \frac{3}{2}\left(Y\Omega_\psi\delta_\psi + \Omega_R\delta_R\right) = 0. \tag{10}$$

For large β^2 the large value of Y accounts for the enhanced attraction between the ψ particles. In addition we observe a generalized damping (or rather antidamping) term $\sim \beta\phi'$ due to the change of the heavy particle mass as ϕ evolves. The perturbations in ρ_ϕ are negligible in view of a unit sound speed. The perturbations δ_R in the radiation fluid follow the standard behaviour and do not grow. Their small amplitude remains at the level inherited from inflation.

Even for a very small initial perturbation δ_ψ the inhomogeneities in the radiation fluid will trigger inhomogeneities in the ψ fluid due to the source term δ_R in the evolution equation for δ_ψ. This typically implies a minimal value for δ_ψ of the same order as δ_R. For the further growth of δ_ψ we can neglect δ_R.

We first discuss the growth rate for the scaling solution (8). Inserting Eq. (9), together with $Y\Omega_\psi \simeq 2/3$ and $\mathcal{H}' = -\mathcal{H}$, and neglecting δ_R, the evolution equation (10) gets rather simple,

$$\delta_\psi'' - \delta_\psi' - \delta_\psi = 0\,. \qquad (11)$$

The solution of this differential equation contains a growing and a decaying mode. At large a, we are left only with the growing piece, namely

$$\delta_\psi = \delta_{\psi,\mathrm{in}}\,(a/a_{\mathrm{in}})^p\,, \qquad p = (1+\sqrt{5})/2 \approx 1.62\,, \quad (12)$$

with $a_{\mathrm{in}} \equiv a(t_{\mathrm{in}})$ the scale factor at the onset of the scaling regime. The ϕ-ψ interactions in Eqs. (4) and (5) translate into a power-law growth of the ψ perturbations during the scaling regime. This is a key difference with respect to non-relativistic matter without the fifth force. For $\beta = 0$ the fluctuations δ_ψ do not grow during radiation domination.

The scaling solution is not essential for the growth of ψ perturbations. For $\Omega_\psi > 1/(3\beta^2)$ the growth is faster than for the scaling solution, while for $\Omega_\psi < 1/(3\beta^2)$ the growth slows down. The growth is also slowed down if ϕ changes slower than the scaling solution (9). Furthermore, the growth rate decreases as the wavelength of fluctuations increases towards the horizon. Many scenarios can be covered by treating p as approximately constant and roughly of the order one. For constant p even small initial inhomogeneities δ_ψ, say $\delta_{\psi,\mathrm{in}} \sim 10^{-5}$, will develop into non-linear inhomogeneities rather rapidly. The number of e-folds for the onset of non-linearity is

$$N_{\mathrm{F}} \equiv \ln\left(\frac{a_{\mathrm{F}}}{a_{\mathrm{in}}}\right) = \frac{1}{p}\ln\left(\frac{\delta_c}{\delta_{\psi,\mathrm{in}}}\right)\,. \qquad (13)$$

with $\delta_\psi(a_{\mathrm{F}}) \equiv \delta_c \sim \mathcal{O}(1)$. This number of e-folds is typically much smaller than the duration of the RD era. For $\delta_{\psi,\mathrm{in}}$ comparable to $\delta_R \approx 10^{-5}$ and $p = (1+\sqrt{5})/2$, it only takes $N_{\mathrm{F}} \approx 7$ e-folds before the fluctuations become non-linear.

We conclude that the fluctuations in the energy density of a non-relativistic fluid always grow non-linear if the following conditions are satisfied: i) the coupling β is large, ii) the fraction Ω_ψ reaches a value of the order β^{-2} not too late in the RD era, such that a value $p \approx 1$ is realized, iii) the scalar field ϕ has a mass smaller than H during the growth period.

An initial Ω_ψ much smaller than $1/(3\beta^2)$ grows until it reaches values of the order $1/(3\beta^2)$ at t_{in}. During this epoch p is small and the growth of fluctuations remains moderate. Once Ω_ψ reaches the scaling regime, $\Omega_\psi \approx 1/(3\beta^2)$, the scaling solution with the fast growth of perturbations (12) becomes a good approximation. For small initial Ω_ψ this simply sets the time t_{in} for the onset of the growth to the time at which Ω_ψ reaches the value $\approx 1/(3\beta^2)$. For more detailed considerations one may employ in Eq. (12) a growth rate $p(a)$ that depends on a.

Black hole formation. When the fluctuations become non-linear one expects the collapse of overdense regions. A collapsing overdense region of the size of the horizon will presumably form a black hole if nothings stops the approximately spherical infall. An alternative would be the formation of highly concentrated lumps. Black holes interact only by gravitational forces. Thus the ψ-particles caught in black holes do no longer feel the fifth force. A similar screened behaviour can be expected for highly concentrated lumps due to strong backreaction effects [17, 18]. For cosmological purposes black holes and screened objects behave similarly and we will no longer make the distinction.

The BH formation subtracts from the energy density ρ_ψ a fraction ρ_{BH} that is converted into BHs. Black holes remain stable even if the ψ particles outside them decay at a later stage. The BH fluid is a non-relativistic fluid contributing to dark matter, with $\rho_{\mathrm{BH}} \sim a^{-3}$ once no new BHs are generated, cf. Fig. 1. A rather small density parameter Ω_{BH} generated early in cosmology will grow $\sim a$ during the long RD period and can reach substantial values. During the BH formation period Ω_{BH} will even grow faster than $\sim a$. We neglect here accretion effects after BH formation which enhance the growth of ρ_{BH}.

There is a simple relation between the time when a black hole forms, as expressed by the scale factor a_{F} at collapse, and the mass of the black hole $M_{\mathrm{BH}}(a_{\mathrm{F}})$. For a rough estimate we assume that all the ψ particles within the horizon at a_{F} form a BH, with mass

$$M_{\mathrm{BH}}(a_{\mathrm{F}}) = \frac{4\pi}{3}\frac{\rho_\psi(a_{\mathrm{F}})}{H^3(a_{\mathrm{F}})} \simeq \frac{4\pi}{3\beta^2}\frac{M_P^2}{H_{\mathrm{F}}}\,. \qquad (14)$$

Here we have employed the value of Ω_ψ according to the scaling solution and $H_{\mathrm{F}} \equiv H(a_{\mathrm{F}})$. Actual masses could be somewhat smaller than the estimate (14) if the infalling mass covers only part of the horizon volume. This does not change orders of magnitude. Expressing M_{BH} in solar mass units $M_\odot$ and taking into account the value of the Hubble parameter at matter radiation equality, H_{eq}, we get

$$\frac{M_{\mathrm{BH}}(a_{\mathrm{F}})}{M_\odot} = \frac{c}{3\beta^2}\frac{H_{\mathrm{eq}}}{H_{\mathrm{F}}}\,, \qquad c \equiv \frac{M_{\mathrm{eq}}}{M_\odot} = 2.7\times 10^{17}\,, \quad (15)$$

with $M_{\mathrm{eq}} = 2GH_{\mathrm{eq}}^{-1}$. Thus BHs with ten solar masses form at an epoch with $H_{\mathrm{F}} \approx 10^{16}\beta^{-2}H_{\mathrm{eq}}$ or equivalently at a redshift $z_{\mathrm{BH}} \approx 10^{11}|\beta|^{-1}$, typically after nucleosynthesis, $z_{\mathrm{NS}} = 10^9$. More massive BHs form even later. We conclude that our mechanism can produce BHs in the mass range observed by the LIGO/VIRGO collaboration only if the mass of the scalar field is below $H_{\mathrm{F}} \approx 2.4 \times 10^{-12}\beta^{-2}$ eV, and if the initial conditions are such that Ω_ψ reaches the scaling solution only near nucleosynthesis.

Let us next estimate the energy density in primordial BHs, ρ_{BH}. Since BH formation proceeds very rapidly once the fluctuations δ_ψ become non-linear, we may assume a complete conversion where all ψ particles end in

BHs. In this limit of instantaneous complete conversion one has $\rho_{\rm BH}(a_{\rm F}) = \rho_\psi(a_{\rm F})$ and $\Omega_{\rm BH}(a_{\rm F}) = \Omega_\psi(a_{\rm F}) \approx 1/(3\beta^2)$. If only a fraction f of the ψ particles ends in black holes our estimate of $\Omega_{\rm BH}$ has to be multiplied by f. After formation $\Omega_{\rm BH}$ grows like non-relativistic matter, such that at the end of the RD epoch one has

$$\Omega_{\rm BH}(a_{\rm eq}) = \frac{a_{\rm eq}}{a_{\rm F}}\Omega_{\rm BH}(a_{\rm F}) = \frac{1}{3\beta^2}\frac{a_{\rm eq}}{a_{\rm F}}\,. \tag{16}$$

For $\Omega_{\rm BH}(a_{\rm eq}) = 1/2$, the BHs constitute all the dark matter in the Universe. An abundance $\Omega_{\rm BH} > 1/2$ would lead to overclosure of the Universe putting bounds on the underlying models with light scalar fields.

Combining Eqs. (15) and (16) we can express the typical mass of the produced BHs in terms of β and $\Omega_{\rm BH}(a_{\rm eq})$,

$$\frac{M_{\rm BH}}{M_\odot} = \frac{c}{3\beta^2}\left(\frac{a_{\rm F}}{a_{\rm eq}}\right)^2 = \frac{c}{27\beta^6\Omega^2_{\rm BH}(a_{\rm eq})}\,. \tag{17}$$

The relations (16) and (17) fix both $M_{\rm BH}/M_\odot$ and $\Omega_{\rm BH}(a_{\rm eq})$ as functions of the parameters β and $a_{\rm F}/a_{\rm eq}$. While β is a model parameter, $a_{\rm F}/a_{\rm eq}$ depends on the initial conditions for ρ_ψ. For fixed β, smaller $\rho_{\psi,{\rm in}}$ leads to larger $a_{\rm F}/a_{\rm eq}$. In particular we may determine the value of β_c for which dark matter is dominated by BHs

$$|\beta_c| = 585\left(\frac{M_{\rm BH}}{M_\odot}\right)^{-1/6} = \left(\frac{2a_{\rm eq}}{3a_{\rm F}}\right)^{1/2}\,. \tag{18}$$

For $\beta^2 \geq 2a_{\rm eq}/(3a_{\rm F}(\beta))$ an additional dark matter component is needed, while models with $\beta^2 \leq 2a_{\rm eq}/(3a_{\rm F}(\beta))$ are excluded.

We finally relate $a_{\rm F}/a_{\rm eq}$ to the initial conditions for Ω_ψ in different scenarios:

1. In our first scenario the ψ particles are produced during the heating and entropy production after inflation. We will denote the end of the heating period by $a_{\rm ht}$. If $\Omega_\psi(a_{\rm ht})$ is of the order $1/(3\beta^2)$ the fluctuations become non-linear at $a_{\rm F}/a_{\rm ht} = \exp(N_{\rm F})$, or equivalently at $a_{\rm F}/a_{\rm eq} = \exp(N_{\rm F} - N_{\rm eq})$ with $N_{\rm eq}$ the number of e-folds between the end of the heating period and matter-radiation equality. If we want to avoid the dimensionless coupling g to be larger than one, β^2 is bounded by M_P^2/m^2. Since m^2 must be larger than $H^2(a_{\rm ht})$, this puts a limit on the end of the heating period,

$$N_{\rm eq} \lesssim \ln\left(\frac{M_P}{H_{\rm eq}}\right)^{2/5} + \frac{N_{\rm F}}{5} \simeq 50 + \frac{N_{\rm F}}{5}\,. \tag{19}$$

The critical value β_c for BH dark matter is very large in this scenario,

$$\beta_c \approx \exp\left[(N_{\rm eq} - N_{\rm F})/2\right]\,, \tag{20}$$

and according to Eq. (17) typical BH masses are tiny as compared to the solar mass. For $\Omega_\psi(a_{\rm ht})$ larger than $1/(3\beta^2)$ the growth of fluctuations will be even faster, with smaller $N_{\rm F}$ and BH masses.

2. For a second scenario we assume $\Omega_\psi(a_{\rm ht}) \ll 1/(3\beta^2)$, while the ψ particles are already decoupled from radiation. In this case the onset of growth is delayed by the time Ω_ψ needs to grow to the order $1/(3\beta^2)$. This replaces $N_{\rm eq} \to N_{\rm eq} + \ln\left(3\beta^2\Omega_\psi(a_{\rm ht})\right)$ in Eqs. (19)-(20). As can be seen from Fig. 1 the relation between ρ_ψ before the scaling solution and $\Omega_{\rm BH}(a_{\rm eq})$ is almost independent of β.

3. In a third scenario the ψ particles are in thermal equilibrium at $a_{\rm ht}$. They decouple from radiation at some time $a_{\rm dc}$. At this time, their abundance can be strongly Boltzmann suppressed, such that $\Omega_\psi(a_{\rm dc})$ can be very small. Effectively, this replaces in Eqs. (19)-(20) $N_{\rm eq} \to N_{\rm eq} + \ln(3\beta^2\Omega_\psi(a_{\rm dc})a_{\rm ht}/a_{\rm dc})$. The BH masses can now be substantially larger, and the LIGO/VIRGO range can be realized for sufficiently small $\Omega_\psi(a_{\rm dc})$.

Conclusions. In this Letter we presented a new mechanism for black hole formation which does not rely on inflationary physics. We argued that primordial black holes could be generated by long-range interactions as those appearing in modified gravity/dark energy scenarios. For illustration purposes we considered a very simple scenario containing just a light scalar field coupled to some fermion field beyond the Standard Model. We showed that for sufficiently large couplings, the system enters a scaling regime in which the fermion energy density *tracks* the background component. During this regime the primordial fermion perturbations become significantly enhanced and can eventually collapse into black holes.

Can the produced black holes contribute to dark matter? According to Ref. [6], there exists an open window for dark matter BHs around 1-1000 $M_\odot$ (see also Refs. [19, 20]). Since this is also the interesting regime to explain the gravitational wave detections, we will focus next on models for which the BH mass distribution peaks at $M_{\rm max} \sim \mathcal{O}(M_\odot)$. The value of $M_{\rm max}$ depends on the precise value of β. For BH dark matter ($\Omega_{\rm BH}(a_{\rm eq}) = 1/2$) it is well approximated by Eq. (18),

$$M_{\rm max} \simeq \left(\frac{585}{\beta}\right)^6 M_\odot\,. \tag{21}$$

A more refined analysis taking into account the entropy production between $a_{\rm F}$ and $a_{\rm eq}$, as well as the distribution of δ_ψ perturbations via a Press-Schechter-type formalism [21], changes the maximum mass (21) by an $\mathcal{O}(1)$ multiplicative factor $\Delta g_s\,(p/(2+p))^{1/p}$, with $\Delta g_s = (g_s(a_{\rm eq})/g_s(a_{\rm F}))^{1/2}$ and g_s the number of entropic degrees of freedom. The mean and the standard deviation of the BH distribution are very close to $M_{\rm max}$, leading to a nearly monochromatic mass spectrum. To obtain masses in the range 1-1000 $M_\odot$, we must have couplings in the range $185 \lesssim \beta \lesssim 585$. For the corresponding very small density parameter in the scaling regime, $\Omega_\psi = 1/(3\beta^2)$, there are no constraints from nucleosynthesis. Black hole formation takes place at an epoch $a_{\mathrm{F}} = 2a_{\mathrm{eq}}/(3\beta^2)$, cf. Eq. (18), corresponding to temperatures of order $\mathcal{O}(\mathrm{MeV})$. The ψ-particles must be stable until this epoch.

Our results are only rough order of magnitude estimates. A more detailed account of the formation and evolution processes is needed. Merging and accretion effects will tend to shift and broaden the BH mass distribution [22, 23]. Extensions and modifications of the model with more than one heavy particles' species as well as decaying particles can be easily constructed. In view of the present uncertainties, a production of black holes in the range $1 - 1000\, M_\odot$ and constituting the whole dark matter component of the Universe seems possible. The necessary ingredients are a scalar field with mass smaller than 10^{-14} eV, "heavy" fields in a suitable abundance, and a mutual effective coupling of the order β_c. The almost massless scalar field could be the cosmon of dynamical dark energy [9, 24], which has at all times a dynamical mass of the order H.

Acknowledgements. We acknowledge support from the DFG through the project TRR33 "The Dark Universe".

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