

## Velocity Dependence from Resonant Self-Interacting Dark Matter

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The dark matter density distribution in small-scale astrophysical objects may indicate that dark matter is self-interacting, while observations from clusters of galaxies suggest that the corresponding cross section depends on the velocity. Using a model-independent approach, we show that resonant self-interacting dark matter can naturally explain such a behavior. In contrast to what is often assumed, this does not require a light mediator. We present explicit realizations of this mechanism and discuss the corresponding astrophysical constraints.

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Dark matter (DM) makes up more than 80% of the matter in the Universe today and played a crucial role in forming stars, galaxies, and hence us. Yet its nature is unknown. Currently, the best pieces of information come from astrophysical observations. *N*-body simulations of collisionless DM predict astrophysical halos with DM density following a universal profile that scales as  $\rho \propto r^{-3}$  in its outskirts but exhibits a central cusp,  $\rho \propto r^{-\beta}$ , with  $\beta \simeq 1$ , referred to as the Navarro-Frenk-White (NFW) profile [1–3]. Nevertheless, many studies show hints of a DM mass deficit in the inner regions of certain halos. Notably, observations indicate that numerous dwarf galaxies [4–6] and some low-surface-brightness spiral galaxies [7–9] have a shallower central DM density, better described by a core of constant density, i.e., by  $\beta \simeq 0$ . This is known as the core-vs-cusp problem. Although it is more pressing in small-scale objects, shallower DM density profiles—with a slope of  $\beta \simeq 0.5$ —have been reported for certain galaxy clusters [10,11]. Moreover, the DM mass deficit also manifests itself in halos that are less dense than what simulations suggest if they host the galaxies that we observe. This is the too-big-to-fail problem, observed for the subhalos of the Milky Way [12], Andromeda [13], and the Local Group [14].

Several explanations for these discrepancies have been discussed in the literature. The systematic uncertainties introduced in deriving DM distributions from observations of luminous objects are one of them. Most importantly, the

motions of HI gas and stars may not be faithful tracers of the DM circular velocity [15–28]. Baryonic processes are another conceivable explanation for the discrepancies, since the aforementioned simulations include only collisionless DM. Solutions along this line include supernova-driven baryonic winds [29–32], DM heating due to star formation [33], and infalling baryonic clumps [34–37] as well as active galactic nuclei or black holes [38]. Nonetheless, there is no consensus on why systematic uncertainties or baryonic processes lead to a seemingly universal mass deficit at various scales.

A more exciting possibility consists of considering DM collisions in the inner regions of astrophysical objects [39]. This is known as self-interacting dark matter (SIDM). *N*-body simulations [40–45] confirm that DM scattering processes indeed reduce the central density of DM halos, providing a solution to both problems [46]. For a recent review, see [51].

The observed mass deficit is more appreciable in small-scale halos, where the DM velocity dispersion is relatively low. Therefore, a self-scattering cross section that decreases with the DM velocity can better fit observations [52], although a constant cross section is certainly not excluded due to the large uncertainties mentioned above. A long-range force induced by a light boson interacting with DM is often invoked to obtain a velocity-dependent cross section [39,53]. Other possibilities that do not involve a light mediator include exothermic inelastic scatterings [54,55] and self-heating DM [56–58].

The essence of this work is to discuss the resonant self-interaction of DM (RSIDM) as another mechanism for achieving the desired velocity dependence of SIDM. Such a resonant behavior was first discussed for DM annihilation in Refs. [59–69] and applied to DM self-scattering in

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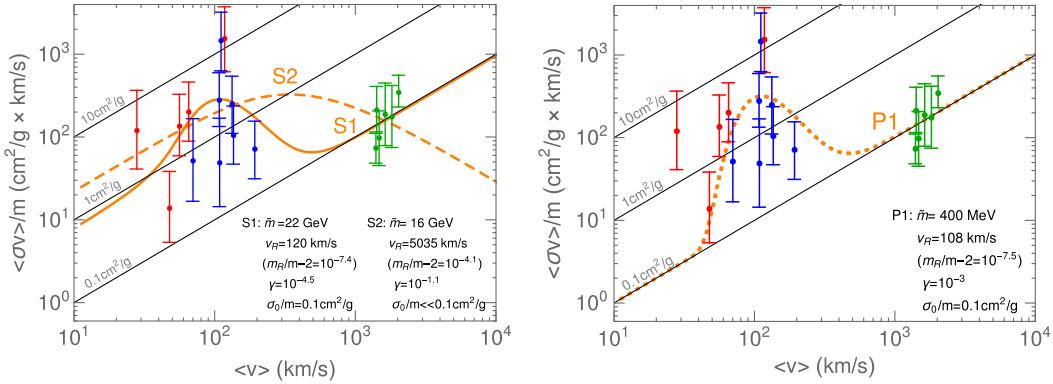


FIG. 1. RSIDM cross section per unit of mass as a function of the velocity. Best-fit curves to the data [52] for  $S$ -wave (left) and  $P$ -wave scatterings (right). The latter is also the best-fit curve for  $L > 1$  after rescaling the mass with Eq. (9). Here  $\tilde{m} = mS^{-1/3}$ . See the text for details.

specific scenarios [70–73]. Nevertheless, the velocity dependence of resonant self-scattering and its general astrophysical consequences have not been explored in detail. In this Letter, we do so in a model-independent way and show that resonant scattering is able to address the observed DM mass deficit at all astrophysical scales. Concrete DM scenarios and indirect searches are discussed later.

*Resonant scattering in DM halos.*—Numerous studies claim that the density distribution of certain DM halos does not follow a NFW profile in the inner region. In the SIDM hypothesis, this is due to DM collisions that thermalize the DM particles in such a region, thereby reducing its average density [39]. Hence, the inner profile is closely related to the velocity-averaged scattering cross section per unit of DM mass,  $\langle\sigma v\rangle/m$ , where [74]

$$\langle\sigma v\rangle = \int_0^{v_{\max}} f(v, v_0) \sigma v dv, \quad f(v, v_0) = \frac{4v^2 e^{-v^2/v_0^2}}{\sqrt{\pi} v_0^3}. \quad (1)$$

Here,  $v$  is the relative velocity, which we assume to follow a Maxwell-Boltzmann distribution truncated at the escape velocity  $v_{\max}$  of the corresponding halo.  $v_0$  is a parameter related to the average relative velocity via  $\langle v \rangle \simeq 2v_0/\sqrt{\pi}$ . Notice that in dwarf galaxies  $\langle v \rangle \sim 20$  km/s, whereas in clusters of galaxies  $\langle v \rangle \sim 2000$  km/s.

A semianalytical method has been proposed in Ref. [52] to infer the value of  $\langle\sigma v\rangle/m$  for a given DM halo from observational data. The method was applied to five clusters from Ref. [11], seven low-surface-brightness spiral galaxies in Ref. [75], and six dwarf galaxies of The HI Nearby Galaxy Survey sample [76] (also see [77,78]). Figure 1 shows their results in green, blue, and red, respectively. The values presented here are for illustrative purpose and should be taken with caution due to the large uncertainties in extracting the cross sections from kinematical data. See, e.g., [79] for a recent study. Nonetheless, at face value, the

figure demonstrates that a cross section independent of the velocity—the ones corresponding to the diagonal lines—can hardly accommodate all points. Notice that the values of  $\sigma/m$  at cluster scales are in agreement with observations from the Bullet Cluster giving  $\sigma/m \lesssim 1.3 \text{ cm}^2/\text{g}$  [80,81], which is one of the strongest constraints on DM self-interactions.

Barring the uncertainties, the figure suggests that the cross section depends on  $\langle v \rangle$ . In this Letter, we propose that this is due to RSIDM. This takes place when there exists an intermediate particle, denoted as  $R$ , so that the total self-scattering cross section can be cast as a sum of a constant piece,  $\sigma_0$ , plus a Breit-Wigner resonance [82]. More explicitly, for nonrelativistic DM,

$$\sigma = \sigma_0 + \frac{4\pi S}{mE(v)} \frac{\Gamma(v)^2/4}{[E(v) - E(v_R)]^2 + \Gamma(v)^2/4}, \quad (2)$$

where the total kinetic energy and symmetry factor read, respectively,

$$E(v) = \frac{1}{2} \frac{m}{2} v^2 \quad \text{and} \quad S = \frac{2J_R + 1}{(2J_{\text{DM}} + 1)^2}. \quad (3)$$

Here,  $J_R$  and  $J_{\text{DM}}$  are the spins of the resonance and the DM particles, respectively.  $m/2$  is the reduced mass. If DM has internal degrees of freedom other than its spin, they must be accounted for in  $S$ . The collision hits the resonance when  $v = v_R$  and, hence,  $E(v_R) = m_R - 2m$ .

In addition, the width in Eq. (2) can be calculated in terms of the resonance self-energy by means of  $\Gamma(v) = \text{Im}\Sigma(v)/m_R$ . This, as well as the denominator in Eq. (2), assumes that the total width is dominated by the process  $R \rightarrow \text{DM DM}$ . Besides that, Eq. (2) is completely general, as it directly follows from unitarity considerations of the scattering matrix [83]. In perturbative theories, the running width can be written as

$$\Gamma(v) = m_R \gamma v^{2L+1}. \quad (4)$$

Here,  $L$  is the orbital angular momentum,  $\Gamma(v_R)$  is the decay rate, and a constant  $\gamma \lesssim O(1)$  characterizes the coupling between the resonance and DM. The factor  $v^{2L+1}$  accounts for the phase space and possible angular momentum suppression. Then we find  $\langle \sigma v \rangle = \sigma_0 \langle v \rangle + 256\pi S \mathcal{I}_L(\gamma, v_R, v_0)/m^2$ , where a dimensionless

$$\mathcal{I}_L(\gamma, v_R, v_0) \equiv \int_0^{v_{\max}} \frac{\gamma^2 f(v, v_0) v^{4L+1} dv}{(v^2 - v_R^2)^2 + 16\gamma^2 v^{2(2L+1)}} \quad (5)$$

determines the nontrivial velocity dependence of the resonant self-scattering. For  $S$ -wave and  $P$ -wave scatterings, we calculate the best-fit parameter sets  $S1$ ,  $S2$ , and  $P1$  based on the inferred data from Ref. [52] and show them in Fig. 1 [84].  $\sigma_0/m$  is fitted with the other parameters for  $S1$  and  $P1$ , while for  $S2$  a negligible  $\sigma_0/m \ll 0.1 \text{ cm}^2/\text{g}$  is taken as a prior. They all lead to  $\chi^2/\text{d.o.f.} \simeq 2$ , in contrast to  $\chi^2/\text{d.o.f.} \simeq 6$  for the fit assuming only a constant cross section (we treat errors as uncorrelated). For  $S1$  and  $P1$ , we show the 95% C.L. contours in Fig. 2. Many comments are in order.

First, we have numerically checked that a precise knowledge of the escape velocity is not necessary for calculating  $\mathcal{I}_L$ . This is because Eq. (5) converges quite fast due to the Boltzmann factor. In fact, exact solutions exist in the limit  $v_{\max} \gg v_0$ , which will be implicitly applied hereafter for simplicity.

Second, to qualitatively understand Figs. 1 and 2, one can use the narrow-width approximation (NWA)

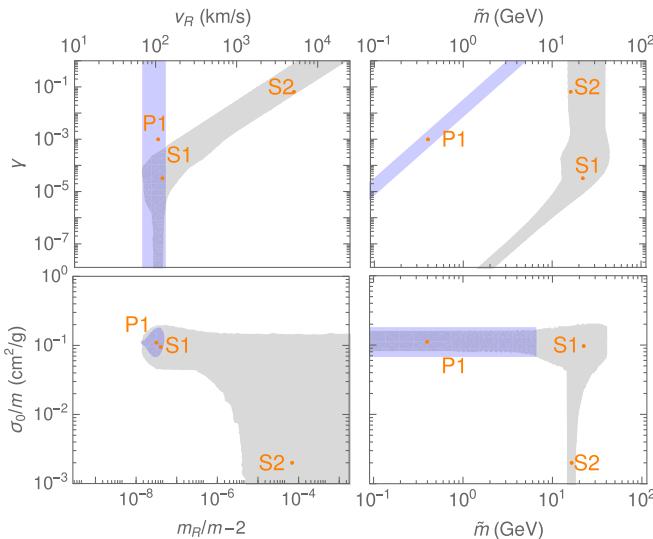


FIG. 2. 95% C.L. contours for  $S$ -wave (gray) and  $P$ -wave (purple) scatterings together with the corresponding parameter sets in Fig. 1. Notice that  $m_R/m - 2 = v_R^2/4$ .

$$\frac{1}{(v^2 - v_R^2)^2 + 16\gamma^2 v^{2(2L+1)}} \rightarrow \frac{\pi \delta(v - v_R)}{8\gamma v_R^{2(L+1)}}. \quad (6)$$

It works very well for  $L \geq 1$ , because  $\gamma^2 v^{2(2L+1)} \ll v^4$ . In this case, we find that  $\mathcal{I}_L(\gamma, v_R, v_0)$  scales as  $\gamma^2 v_0^{4L+1}/v_R^4$  at  $v_0 \ll v_R$  and as  $\gamma^2 v_0^{4L-3}$  at  $v_0 \gg v_R$ . In both regions,  $\mathcal{I}_L$  cannot be much larger than one. Therefore, the resonant effect is negligible except for the intermediate region, where the NWA captures the velocity dependence as

$$\left. \langle \sigma v \rangle \right|_{\text{NWA}} = \frac{\sigma_0 \langle v \rangle}{m} + \frac{128S\pi^{3/2}\gamma v_R^{2L+1}}{m^3 v_0^3} e^{-v_R^2/v_0^2}. \quad (7)$$

Notice that the peak lies at  $v_0 \sim v_R$  as illustrated by  $P1$  in Fig. 1. The corresponding line actually applies to any  $L \geq 1$ , because the dependence on  $L$  can be absorbed by rescaling  $m$ . Using Eq. (7), we find that the best-fit parameters at 95% C.L. for  $L \geq 1$  are given by

$$v_R = [108^{+28}_{-43}] \text{ km/s}, \quad \sigma_0/m = [0.11^{+0.10}_{-0.05}] \text{ cm}^2/\text{g}, \quad (8)$$

$$\tilde{m} = [400^{+120}_{-90}] \text{ MeV} \left( \frac{\gamma}{10^{-3}} \right)^{1/3} \left( \frac{v_R}{[3 \times 10^5] \text{ km/s}} \right)^{2(L-1)/3}. \quad (9)$$

Such values for the velocity correspond to  $m_R/m - 2 - 10^{-7}$ . The regions where all this applies are shown in Fig. 2. For  $P$ -wave scattering, demanding  $\gamma \lesssim 1$  leads to  $\tilde{m} \equiv mS^{-1/3} \lesssim 5 \text{ GeV}$ . Moreover, a perturbative  $\sigma_0/m$  around  $0.1 \text{ cm}^2/\text{g}$  requires sub-GeV DM masses unless  $S \gg 1$ . Interestingly,  $P1$  predicts  $\sigma/m \sim 0.1 \text{ cm}^2/\text{g}$  at  $\langle v \rangle \ll 100 \text{ km/s}$ . In fact, scatterings with  $L \geq 1$  can realize small cross sections at very low velocities. Hence, the recent claim based on Draco observations [28] is consistent with RSIDM.

As long as  $v_R \gtrsim 4\gamma$ , the NWA also applies for  $S$ -wave scattering. For  $v_R \ll 4\gamma$ ,  $\mathcal{I}_L$  is proportional to  $v_0$  (to  $1/v_0$ ) below (above)  $v_{\text{peak}} \sim v_R^2/(4\gamma) \ll v_R$ , because such large values of  $\gamma$  broaden the resonance.  $S1$  and  $S2$  illustrate the narrow and the broad width cases, respectively.

In conclusion, resonant scattering is able to address the observed DM mass deficit at all astrophysical scales.

**RSIDM models.**—Below, we illustrate the previous model-independent results in concrete RSIDM scenarios. We first introduce a Lagrangian specifying the coupling of the DM to the resonance (see Table I) and calculate the cross section and the self-energy. We subsequently corroborate that they can be cast as Eqs. (2) and (4) show. The scenarios are as follows.

**(I) Fermionic DM with a pseudoscalar mediator.**—The scattering process is  $S$ -wave while  $\sigma_0 \simeq 0$ . The corresponding best fit is thus  $S2$ . Notice that a light pseudoscalar mediator does not lead to SIDM, because it induces a

TABLE I. Benchmark RSIDM models.

Scenario	Interaction Lagrangian	$L$	$J_{\text{DM}}$	$J_R^P$	$S$	$\gamma$
I	$gR\bar{D}\bar{M}\gamma^5DM$	0	$\frac{1}{2}$	$0^-$	$\frac{1}{4}$	$g^2/32\pi$
IIa	$gR\bar{D}M^i\bar{D}M^i$	0	0	$0^+$	$\frac{1}{3}$	$g^2/16\pi m_R^2$
IIb	$g\epsilon_{ijk}R_\mu^i\bar{D}M^j\partial^\mu DM^k$	1	0	$1^-$	1	$g^2/384\pi$
III	$(1/\Lambda)R_{\mu\nu}\bar{T}_{\text{DM}}^{\mu\nu}$	2	0	$2^+$	5	$m_R^2/30720\pi\Lambda^2$

suppressed Yukawa potential (see, e.g., [85]). Because of this and because it leads to velocity-suppressed direct-detection rates, this candidate is phenomenologically interesting.

(II) *Dark mesons*.—In QCD-like theories, DM can be a dark pion. Analogous to real pions, it can be a triplet  $DM^i$ , with  $i = 1, 2, 3$ . If  $R$  is a dark  $\sigma$  resonance (IIa), the scattering takes place via the  $S$  wave, where we expect GeV DM and  $\sigma_0/m \ll 100 \text{ cm}^2/\text{g}$ . The best fit is thus  $S2$ . If  $R$  is a dark  $\rho$  resonance (IIb), the scattering is  $P$ -wave suppressed. The constant piece of the cross section is given by  $\sigma_0 \sim \pi\gamma^2/m^2$  in the perturbation theory, but it is plausible that there are other contributions. We therefore leave  $\sigma_0$  as a free parameter. The corresponding best-fit curve is  $P1$ . We expect  $m \sim 400 \text{ MeV}$  in this case. In the same fashion, minimal QCD-like theories can also lead to spin-1 DM [86]. In all cases, DM can be produced by means of the strongly interacting massive particle (SIMP) [87–109] and the freeze-in [110–112] mechanisms.

(III) *Tensor resonances*.—They also arise in strongly coupled theories. Despite the potential complications of such theories, the generality of our approach allows us to describe the scattering induced by a spin-2 resonance  $R_{\mu\nu}$  [113]. If this couples to the DM energy-momentum tensor with a cutoff scale  $\Lambda$ , and taking scalar DM as an example, we find that the corresponding Feynman rules [114] indeed lead to a  $D$ -wave cross section given by Eq. (2). For  $m \sim 10^{-3}\Lambda$ , we obtain keV DM with  $\gamma \sim 10^{-13}$ . The corresponding best fit is given by  $P1$  in Fig. 1 after rescaling the mass by means of Eq. (9).

*Annihilation vs scattering*.—It is not necessary that the DM annihilates, as, e.g., in models of asymmetric DM. Nonetheless, if the resonance decays into a pair of standard model (SM) particles  $f\bar{f}$ , in analogy to Eq. (2), the resonant DM annihilation into  $f\bar{f}$  has a cross section

$$\sigma_{\text{anni}} \simeq \frac{4\pi S}{mE(v)} \frac{\Gamma(v)m_R\gamma_f/4}{[E(v) - E(v_R)]^2 + \Gamma(v)^2/4}, \quad (10)$$

where  $m_R\gamma_f$  is the decay width for  $R \rightarrow f\bar{f}$ . As above, we assume that the resonance dominantly decays to a pair of DM particles and, thus, that the contribution of  $f$  to the imaginary part of the resonance self-energy,  $m_R^2\gamma_f$ , is subleading. This is different from Ref. [70], in which the resonance dominantly decays into visible particles.

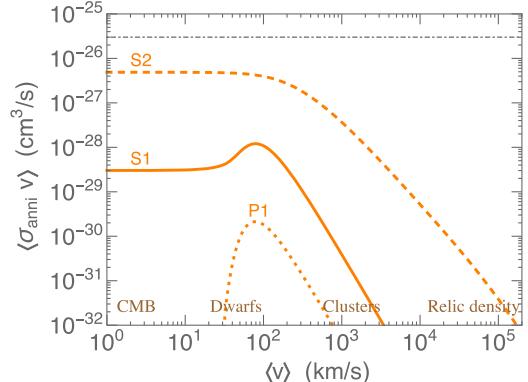


FIG. 3. Annihilation cross section into a pair of charged fermions for the parameter sets in Fig. 1, assuming a branching ratio  $\gamma_f/(\gamma v_R^{2L+1}) = 10^{-13}$ . The horizontal line gives the standard freeze-out benchmark.

As expected for annihilations (but not for elastic scatterings),  $\sigma_{\text{anni}} v \propto v^{2L}$  as long as  $v \ll v_R$ . Furthermore, for the cases where NWA applies,  $\langle\sigma_{\text{anni}} v\rangle|_{\text{peak}} \sim 32\pi^2 S\gamma_f/(m^2 v_R^3)$ . In contrast, for broad  $S$ -wave resonances such as  $S2$ , where  $v_{\text{peak}} \ll v_R$ ,  $\langle\sigma_{\text{anni}} v\rangle|_{\text{peak}}$  gets enhanced by another factor  $(v_R/v_{\text{peak}})^{2L+1}$ .

The coupling to light charged particles is mostly constrained by Fermi-LAT observations of local satellites [115,116] and the Planck data on the cosmic microwave background (CMB)[117,118]. For instance, the corresponding Fermi-LAT upper limit on  $\langle\sigma v_{\text{anni}}\rangle$  for GeV DM is of the order of  $10^{-26} \text{ cm}^3/\text{s}$ . For  $S2$ , this leads to an upper limit on the branching ratio,  $\gamma_f/(\gamma v_R)$ , of about  $10^{-13}$ – $10^{-12}$ . This bound is much stronger than that of  $S1$  and  $P1$ , due to the enhancement factor mentioned above. Motivated by this, we conservatively fix  $\gamma_f/(\gamma v_R^{2L+1}) = 10^{-13}$  and calculate the annihilation cross section as a function of  $\langle v \rangle$  for the same parameter sets in Fig. 1. The result is shown in Fig. 3. Therefore, the resonance can couple only feebly to light charged particles, which is why the SIDM candidates with thermal freeze-out from Ref. [72] are excluded. Of course, this is model dependent. For instance, if the resonance couples only to neutrinos, the bound on  $\langle\sigma_{\text{anni}} v\rangle$  becomes much weaker, and larger  $\gamma_f/\gamma$  are thus allowed.

Furthermore, the strong velocity dependence of  $\langle\sigma v_{\text{anni}}\rangle$  suggests that the usual freeze-out can hardly work, as for SIDM with light mediators decaying into visible particles [119–122]. Nevertheless, the DM abundance might arise from other SIDM production mechanisms [112]. Indeed, for the  $S$ -wave case, producing the DM abundance with small couplings is possible via freeze-in [110,111] or 4-to-2 annihilations [97], where a scalar (vector) resonance can feebly mix with the Higgs (SM gauge bosons). See [123,124] for reviews.

*Discussion*.—We advocate the resonant scattering as a possible SIDM realization with a velocity-dependent

scattering cross section. Instead of a light mediator, this RSIDM scenario requires a near-threshold resonance with  $m_R/m - 2$  ranging from  $\sim 10^{-7}$  for narrow resonances to  $10^{-2}$  for  $S$ -wave scattering with broad widths. Such resonances exist in nature. As an example,  $\alpha$  particles resonantly scatter by means of  $^8\text{Be}$  in exactly the same way as described above. In fact, these processes were the main subject of the original article by Breit and Wigner [83], and they may as well occur in the DM sector. Actually, dark nucleons as SIDM have been studied in Refs. [73,125]. Furthermore, lattice studies suggest that QCD-like theories of DM might possess such states [126].

*Conclusions.*—We find that this RSIDM hypothesis can certainly address the core-vs-cusp and the too-to-big-fail problems while still being in agreement with cluster observations. We have also discussed indirect detection signatures, which are nevertheless model dependent. Additionally, we would like to emphasize that usual SIMPs—which are often said to be disfavored because their scattering cross section does not vary with velocity—can easily accommodate the mechanism proposed here.

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