Dark matter filtering-out effect during a first-order phase transition

Dongjin Chway,1,* Tae Hyun Jung,1,2,† and Chang Sub Shin1,‡

1Center for Theoretical Physics of the Universe, Institute for Basic Science (IBS), Daejeon, 34051, Korea
2Department of Physics, Florida State University, Tallahassee, FL 32306, USA

If the mass of dark matter is originated from a symmetry breaking and the corresponding phase transition is first-order, the bubble walls during the phase transition can filter out dark matter particles. Only particles which have enough momentum to overcome their masses inside the bubbles can penetrate through the bubble walls. Consequently, the dark matter number density after the phase transition has a suppression factor of \( \exp(-\frac{M_\chi}{2\gamma T}) \) where \( M_\chi \) is the dark matter mass, and \( \gamma \) and \( T \) are, respectively, the Lorentz gamma factor and temperature of the incoming fluid in the bubble wall rest frame. We show that the filtering-out process provides sizable pressure on the bubble wall to obtain \( \gamma T \) small enough, which makes the suppression factor consistent with the observed dark matter relic abundance for a wide range of dark matter masses up to the Planck scale.

Introduction

The thermal freeze-out mechanism of dark matter (DM) has been the mainstream of understanding the DM abundance [1]. As temperature drops, when the DM annihilation rate becomes smaller than the Hubble expansion rate, the DM number changing interactions are no more in thermal equilibrium and its comoving number density is frozen. To explain the observed DM relic density, DM particles need a sizable annihilation rate to lighter particles, roughly as large as the electroweak interaction rate. We call such hypothetical DM particles as weakly interacting massive particles (WIMP).

Motivated by the WIMP paradigm, there have been a lot of experimental studies to reveal its particle properties. Especially, direct detection experiments using large-volume detectors have increased their sensitivity enormously for the last decades. Unfortunately, we have not yet obtained the convincing signal in direct detection experiments so far. The absence of a direct detection signal provides strong constraints on the simple WIMP dark matter models with masses from GeV to TeV scale [2].

Even if we can take refuge in heavy WIMP scenarios, there is the strong upper bound on DM mass within the freeze-out mechanism. The upper bound comes from the fact that the required annihilation cross section increases as the DM mass increases, and the cross section should not exceed the unitarity bound. The unitarity bound implies that the WIMP mass to be less than around 100 TeV [3, 4].

Therefore if the DM mass is heavier than 100 TeV, there should be an additional process to provide the correct relic density. Along this direction, the pioneering works [5–9] studied the role of the early matter domination and inflation periods to obtain the correct heavy DM relics. The freeze-in thermal production [10], and the series of co-scattering processes [11] are also used to make the thermal heavy WIMP DM scenario viable. One of the easiest ways to overcome the problem is the entropy injection to the SM after freezing out DM, which could be originated from a supercooled first-order phase transition of the Universe [12, 13].

In this paper, we show that a first-order phase transition has another intrinsic effect which reduces the DM density compared to that from the naive freeze-out process; the bubble wall of the phase transition can filter out DM particles (see Fig.1). We call this mechanism as filtering-out mechanism.

Filtering Effect

Let us consider that the DM mass is originated from a symmetry breaking and the corresponding phase transition is first-order. During the phase transition, symmetry breaking bubbles are nucleated and expand. The DM mass is nonzero inside the bubbles while it is zero outside the bubbles. Therefore, only DM particles that have enough energy can penetrate through the bubble walls.

In the wall rest frame, the number of particles penetrating through the bubble wall per area \( \Delta A \) and a time interval \( \Delta t \) can be estimated by

\[
\frac{\Delta N_m}{\Delta A \Delta t} = \frac{g_\chi}{(2\pi)^3} \int d^3\tilde{p} \int_{r_0}^{r_{wall}} dr \mathcal{T}(\tilde{p}) \Theta(-p_z) f(\tilde{p}, \vec{x}) ,
\]

(1)

where \( r_0 \) denotes the bubble radius, \( \mathcal{T}(\tilde{p}) \) is the transmission rate, \( f(\tilde{p}, \vec{x}) \) is a distribution function of particles at position \( \vec{x} \) and momentum \( \tilde{p} \), \( g_\chi \) is the particle degrees of freedom, and \( \Theta(x) \) is the unit step function. For simplicity, we take classical transmission rate \( \mathcal{T}(\tilde{p}) \approx \Theta(-p_z - M_\chi) \) where \( M_\chi \) is the DM mass inside the bubble.

If the distribution follows Bose-Einstein or Fermi-Dirac distribution

\[
f_{eq}(\tilde{p}, \vec{v}, T) = \frac{1}{e^{\gamma (E - \tilde{p} \cdot \vec{v}) / T} \pm 1} ,
\]

(2)

then the particle flux \( J_{wall} = dN/(dAdt) \) in the bubble...
FIG. 1. A schematic of the filtering-out mechanism. Most of DM particles $\chi$ cannot penetrate through the bubble wall if, in the wall rest frame, momenta of particles outside the bubbles are not high enough to overcome the mass that DM would obtain inside the bubbles. DM particles annihilate into other light particles which can more freely enter the bubble.

The flux density in the bubble can be obtained by

$$J_{\text{wall}} = \frac{g_s T^3}{2 (2\pi)^2 \gamma (1 - \tilde{v})} \int_{x_*}^{\infty} \frac{dx}{e^{x} - 1} (x^2 - x^2),$$

where $x_* = \tilde{\gamma}(1 - \tilde{v}) M_\chi / T$ with $\tilde{\gamma} = 1/\sqrt{1 - \tilde{v}^2}$ being the Lorentz gamma factor of the relative velocity $\tilde{v}$ of fluid bulk motion with respect to the wall. With Boltzmann distribution, dropping $\pm 1$ in the denominator, we obtain the flux

$$J_{\text{wall}} \simeq - g_s T^3 \left( \frac{\tilde{\gamma}(1 - \tilde{v}) M_\chi / T + 1}{4 \pi^2 \gamma^3 (1 - \tilde{v})^2} \right) e^{-\tilde{\gamma}(1 - \tilde{v}) M_\chi / (1 - \tilde{v})^2}.$$

For Bose-Einstein or Fermi-Dirac distribution, we obtain only a factor of $O(1)$ change. Note that even for $M_\chi \gg T$, the exponent is quite sensitive to $\tilde{v}$. It becomes $-M_\chi / 2 \gamma T$ ($-M_\chi / T$) as $\tilde{v} \to 1$ ($\tilde{v} \to 0$). Including the case of the supercooled Universe, a natural environment for dark matter filtering process, we allow all the possible range of $\tilde{v}$. There is also the outgoing flux which corresponds to $-J_{\text{wall}}(\tilde{v} \to -\tilde{v})$. Because of the exponential factor, for the sizable $\tilde{v}$, this can be safely ignored compared to the ingoing DM flux, while it becomes important when $\tilde{v} \ll 1$.

Under the assumption that DM particles inside the bubble are thermally decoupled, the average number density in the bubble can be obtained by

$$n_\chi = -\frac{J_{\text{wall}}}{\gamma w \xi w},$$

where $\xi w$ is the bubble wall velocity and $\gamma w$ is its gamma factor. When $\tilde{\gamma} \gg M_\chi / T$ and $\xi w \simeq \tilde{v}$, Eq. (5) approaches $g_s T^3 / \pi^2$ which is the equilibrium number density for Boltzmann distribution outside the bubble. It means that the bubble wall does not filter out DM particles at all in this limit.

Plasma which is farther than the mean free path from the bubble wall can be described by a perfect fluid distribution with $f_{\text{eq}}(\vec{p}, \vec{v}, T)$ in Eq. (2) where $\vec{v}$ and $T$ should be treated as functions of the distance from the center of the bubble. As described in Ref. [14], the conservation of energy and momentum provides the profiles of $\vec{v}$ and $T$ once two effective parameters are given: supercooling order parameter $\alpha_n$ and the bubble wall velocity $\xi w$.

The $\alpha_n$ describes the potential difference between false and true vacua,

$$\alpha_n = \Delta V / (a_n T_n^4),$$

where $\Delta V$ is the potential energy difference at zero temperature, and the sub-index $n$ denotes a quantity at the bubble nucleation temperature $T_n$. Here, $a$ is the effective light degrees of freedom at a given temperature

$$a = \frac{\pi^2}{30} \sum_{i \text{ light}} \left( N_i^b + \frac{7}{8} N_i^f \right),$$

where $N_i^b$ and $N_i^f$ are bosonic and fermionic field degrees of freedom.

With nonzero $\Delta V$ and infinitesimal bubble wall width (compared to the macroscopic bubble size), the fluid profile becomes discontinuous at the bubble wall, which can be determined from the conservation of the energy and momentum. We denote velocity relative to the wall and temperature of the incoming (transmitted) fluid at the bubble wall by $v_+$ ($v_-$) and $T_+$ ($T_-$), following the notation of Ref. [14].

The distribution close the wall, less than the mean free path away, should be determined, in principle, by solving the Boltzmann equation. In this paper, however, instead of attempting to solve it, we make assumptions on the out-of-equilibrium distribution to be approximately described by

$$f(\vec{p}) = \begin{cases} f_{\text{eq}}(\vec{p}, -\vec{v}, T) + f_{\text{eq}}(\vec{p}, \vec{v}, T) & \text{if } p_r < M_\chi / 2 \\ f_{\text{eq}}(\vec{p}, -\vec{v}, T) & \text{if } p_r > M_\chi \end{cases}$$

In the main discussion, we further set

$$\tilde{v} = v_+ \text{ and } T = T_+. \quad (9)$$

This corresponds to the situation when the self interaction of the reflected particle species is negligible. If the incoming DM particles which will be reflected are not affected by the DM particles which was reflected, the distribution of incoming DM particles will not change much. In such a case, it is natural to set $\tilde{v} = v_+$.

We can calculate $\xi w$, $\tilde{v}$ and $T$ from the equilibrium condition $\Delta V = P$. If a particle $i$ gains its mass $M_i$ inside the bubble and $\tilde{\gamma} T \lesssim 0.2 M_i$, most of the $i$ particles in the flux fail to penetrate and exert pressure on the wall.1

---

1 If $\tilde{\gamma} T$ is larger than $M_\chi$, the pressure is $P = \frac{\gamma w}{\alpha} M_\chi^2 T^2$ where $\alpha_1$ is 1 for boson and 1/2 for fermion [15]. Its next leading order is discussed in Ref. [16].
The pressure can be represented as
\[ P = \frac{d}{3} a(1 + \dot{v})^2 \gamma^2 T^4, \] (10)
defining the effective ratio of the degrees of freedom by
\[ d \equiv \frac{1}{a} \left[ \frac{\pi^2}{30} \sum_{\text{all } M_i > \gamma T} \left( N_i^b + \frac{7}{8} N_i^f \right) \right]. \] (11)

This is nonzero in our scenario because we have at least one heavy particle: DM.

**Bubble collisions** When the bubbles collide with each other, bulk kinetic energy of plasma can be converted into the local thermal energy. During this process, locally very high-temperature plasma can be formed for a while and may produce dark matter particles when the phase transition is supercooled.\(^2\) When supercooling occurs, the bulk motion of fluid profile can be described by a detonation profile or a hybrid profile [14]. In detonation profiles, hot plasma follows the bubble wall, and the bubble wall hits the plasma that is at rest. Thus, \(v_-\) and \(T_-\) are important at the collision. In hybrid profiles, hot plasma is additionally involved in front of the bubble wall propagation, and so \(v_+\) and \(T_+\) should be taken.

This thermalization process of colliding fluids should not be ignored because the mean free path can be short compared to the thickness of the plasma profile. The mean free-path is roughly \(l \sim \gamma/(g^4 T)\) where \(g^4\) represents interaction strength. On the other hand, the thickness of the plasma profile is approximately \(R/\alpha_n\) for large \(\alpha_n\) where \(R\) is the average radius of the bubbles at the percolation time. In this paper, to give a conservative estimation, we assume that \(R\) is cosmologically large.

At the collision, the plasma energy density \(\gamma_{\text{pl}}^2 T^4\) can be converted into thermal energy. Here, \(\gamma_{\text{pl}}\) is the Lorentz gamma factor of the plasma fluid in the bubble center rest frame. Since \(\gamma_{\text{pl}} \sim \gamma_w \propto \alpha_n^{1/2}\) from \(\Delta V = P\) with Eq. (10) and \(T_- \sim \alpha_n^{1/4} T_n\) for detonations and \(T_+ \sim \alpha_n^{1/4} T_n\) for hybrids, the maximal temperature of the local fluid can be as large as
\[ T_{\text{max}} \sim \alpha_n^{1/2} T_n, \] (12)
for both detonation and hybrid profiles. We restrict our scenario to satisfy \(T_{\text{max}} < T_{\text{fo}}\) to prevent DM from being additionally produced where \(T_{\text{fo}}\) is the DM freeze-out temperature. It is a conservative bound since the additional DM production in this process would not be efficient depending on the interaction strength.

\(^2\) When \(\Delta V\) is large, collisions of scalar profiles generate scalar field oscillations without producing a significant amount of heavy DM [17]. If the scalar field is derivatively coupled to light particles, the oscillation energy decays into the light particles.

In hybrid profiles, the shock-wave plasma collides first, and high-temperature plasma can be formed between bubble walls. The DM production is efficient since DM is still massless between the bubble walls. If there is enough time between the collision of shock waves and the collision of scalar profiles, a huge amount of DM particles can be produced and smear into the bubbles. The time scale between shock wave collision and bubble wall collision can be estimated by \(\Delta t \sim R(\xi_{sh} - \xi_w) \approx R/\alpha_n\) where \(\xi_{sh}\) is the shock-front velocity. In order to be conservative in our estimation, we have assumed that \(R/\alpha_n < \gamma/(g^4 T)\), so we conclude that DM can be easily overproduced during the bubble collision if the fluid profile is hybrid. Thus, we exclude the hybrid scenario from our consideration.

**Results** We define an effective parameter \(\lambda_{\text{eff}}\) to represent how small \(\Delta V\) is,
\[ \lambda_{\text{eff}} = \frac{\Delta V}{M_X^2}, \] (13)
compared to the DM mass \(M_X\) inside the bubbles. In terms of \(T_n\) and \(\lambda_{\text{eff}}\), we can rewrite \(M_X/T_n \approx 2.4(\alpha_n/\lambda_{\text{eff}})^{1/4}\). Thus, the DM relic density can be written in terms of \(\alpha_n, \lambda_{\text{eff}}, d\) and \(T_n\).

When there is no supercooling \((\alpha_n \lesssim 0.1)\), the wall velocity remains small. Thus, in the approximation with \(\gamma \simeq 1\) and \(T \simeq T_n\), the exponential suppression factor in Eq. (4) and (5) becomes \(\exp[-2.4(\alpha_n/\lambda_{\text{eff}})^{1/4}]\). Numerically, \(\Omega h^2 \lesssim 0.1\) requires that
\[ \frac{\lambda_{\text{eff}}}{\alpha_n} \simeq \left(10.2 + 0.45 \log \frac{T_n}{\text{GeV}} \right)^{-4} \] for \(\alpha_n \lesssim 0.1\). (14)

However, we find that, as discussed in the next section, making small \(\alpha_n\) with small \(\lambda_{\text{eff}}\) seems difficult to be realized unless some strong interactions are introduced.

For a supercooled first-order phase transition \((\alpha_n \gtrsim 1)\), the situation becomes more complicated. The phase transition involves not only filtering-out effect, but also the bubble collision effect and reheating process.

With a detonation fluid profile, we show numerical value of \(T_n\) required for the correct DM relic density for \(d = 0.05\) in Fig. 2. In this plot, we set \(a_n = 100 \times \pi^2/30\) and DM degrees of freedom 2. From the condition that DM particles outside the bubble should annihilate efficiently, we have \(\Gamma_{\text{ann}} \ll H\) where \(\Gamma_{\text{ann}}\) is annihilation rate and \(H\) is Hubble rate. Therefore, it needs to satisfy \(T_n \ll M_{\text{Pl}}\). The purple shaded region describes \(T_{\text{max}} < T_{\text{fo}}\) where one might have to consider DM produced during the bubble collision as discussed in the previous section. For each \(M_X\), we obtained \(T_{\text{fo}}\) numerically as in Ref. [18], using \(\langle se\rangle \simeq \alpha^2/M_X^2\) with \(\alpha \simeq 1\) to be conservative.

It is noteworthy that, in the figure, the bottom-left corner and the top-right corner can be described by dif-
FIG. 2. $T_n$ required for the observed DM relic is depicted in the $\alpha_n$-$\lambda_{\text{eff}}$ plane with the detonation profile. Black solid lines indicate the DM mass $M_\chi$. The purple region is where the DM particles are produced too much during the bubble collisions.

Different schemes. Since $\tilde{\gamma} = \gamma/w \sim \alpha_n^{1/2}$ and $T = T_n$, in the bottom-left region, the exponent becomes $-M_\chi/\tilde{\gamma}T \sim -1/(\lambda_{\text{eff}}\alpha_n)^{1/4}$ which makes the filtering effect dominant. The degeneracy can be seen along the line $\lambda_{\text{eff}}\alpha_n \sim 10^{-9}$. In order for heavier DM to satisfy the observed relic, we need smaller $\lambda_{\text{eff}}\alpha_n$ which enters with the exponential factor, roughly by $\exp(-1/(\lambda_{\text{eff}}\alpha_n)^{1/4})$. Therefore, small changes in $\alpha_n$ and $\lambda_{\text{eff}}$ can allow the DM arbitrarily massive even up to the Planck scale, as can be seen in Fig. 2.

On the other hand, in the top-right corner, most of the DM particles just enter the bubble wall because $\tilde{\gamma} \gg M_\chi/T$. In this case, the reheating process provides the suppression factor $(T_n/T_\text{rh})^3$ required for the observed DM relic density [12, 13]. However, in this region, DM particles might be additionally produced during the bubble collision since $T_{\text{max}} > T_\text{rh}$, and so more careful study is required.

If DM self interaction is strong enough, Eq. (9) could not be valid. In such a case, we can consider two simple possibilities under the condition that there exists a terminal velocity of the bubble wall: 1) $\tilde{v} = v_\gamma$ and $T = T_\gamma$, and 2) $\tilde{v} = 0$ and $T \sim T_\text{rh}$. Both cases result in the exponential suppression factor $\exp[-\lambda_{\text{eff}}^{1/4}]$ without $\alpha_n$ dependence. Here, the dilution factor disappears since $T \sim T_\text{rh}$. We would like to emphasize that it is important to solve full Boltzmann equations around the bubble wall especially for a strong DM self interaction. In the worst case, because of chaotic motions of DM particles in front of the bubble wall, pressure might not be enough for the bubble wall terminal velocity to exist.

Discussion In this paper, we have investigated the possibility that the DM relic abundance is determined by the filtering-out effect of the bubble wall during a first-order phase transition. We have shown that the DM number density after phase transition is suppressed by $\exp(-M_\chi/2\gamma T)$. Unlike the freeze-out mechanism, our scenario does not have any theoretical lower bound on DM number density so that the DM mass can be as large as the Planck scale. In terms of effective parameters $\alpha_n$ and $\lambda_{\text{eff}}$, we find a phenomenologically reasonable region for the observed DM relic density. For the detonation profile, $\lambda_{\text{eff}}\alpha_n$ should be order of $10^{-9}$.

An intrinsic observable of this mechanism is a gravitational wave signature since a strong first-order phase transition is required, $T_n \ll M_\chi$. Gravitational wave produced in a first-order phase transition has been widely studied in various contexts [19–26]. The signal peak frequency is, roughly, $1/R$ multiplied by a redshift factor where $R$ is the bubble radius at the bubble collision. To estimate the signal strength, we need to specify the model, but it can be arbitrary at this moment. If we have more observational information from future gravitational wave detectors [26–32], we will be able to narrow down $T_n$, $M_\chi$ and $\Delta V$ required.

Indirect detection is still a good possibility to check our scenario. The filtering-out mechanism requires DM particles to annihilate into light particles, efficiently, when they were massless. If not, they are accumulated outside the bubble, heated, and finally penetrate through the bubble wall when the temperature becomes high enough. In order to prevent this heating process in the symmetric phase, they should annihilate into light particles that can freely enter the bubble. Particles in the standard model such as neutrinos and photons are good candidates although it is not necessary. Excess in the high energetic cosmic rays is one of the footprints that our mechanism can have.

As a final remark in the model building aspect, we note that the scalar potential should contain at least two different mass scales. Let us first consider a maxican hat potential $V = -m^2\Phi^2 + \lambda\Phi^4$ which has only one massive parameter $m$. Given a Yukawa coupling $y_\chi$ between $\Phi$ and the dark matter, $\lambda_{\text{eff}} \sim \lambda/y_\chi^4$, $T_n \sim m/y_\chi$, and we obtain $\alpha_n \sim 10^{-4}y_\chi^2/\lambda$. For $\alpha_n \lesssim 0.1$, the model cannot satisfy Eq. (14). Even if we consider a large $\alpha_n$, we obtain $\lambda_{\text{eff}}\alpha_n \sim 10^{-4}$ which is much bigger than the value for the correct DM relic, $O(10^{-9})$.

One of the working examples to provide multiple scales is the supersymmetric (SUSY) axion model in gauge mediation with a messenger scale $M \ll M_p$. The shapes of the scalar potential for the saxion (the superpartner of the axion) field are quite different between two regions $\Phi < M$ and $\Phi > M$. In the field range $\Phi < M$, soft SUSY breaking mass terms are generated by gauge mediation so that $V \sim -m_a^2\Phi^2$, while for $\Phi > M$ its effect is quite suppressed and the potential becomes $-m_a^2M^2(\ln \Phi/M)^n + m_a^2/2\Phi^2$. [33–41]. The gravitino...
mass $m_{3/2}$ is much smaller than $m_s$, so the vacuum value of the saxion is evaluated as $\langle \Phi \rangle \sim M m_{3/2}/m_{3/2} \gg M$. In that case, $T_n \simeq m_s$, $\Delta V \sim m_s^2 M^2$, and one can easily obtain $\lambda_{\text{eff}} \alpha_n \sim (m_{3/2}/m_s)^4 = O(10^{-9})$. With this potential, Eq. (14) with a small $\alpha$ requires $m_s = O(M)$, which is very unlikely in gauge mediation scenario. We leave detailed studies within specific models as future works.

Acknowledgement This work was supported by IBS under the project code, IBS-R018-D1. T.H.J. is supported by the US Department of Energy grant DE-SC0010102 and Prof. Kohsaku Tobioka’s startup fund at Florida State University (Project id: 084011-550-042584).

Note added While this paper was under completion, the Ref. [43] appeared on arXiv, which is based on the same idea but different implementations.

---

* djchway@gmail.com
† thjung0720@gmail.com
coshin@ibs.re.kr

[1] B. W. Lee and S. Weinberg, “Cosmological Lower Bound on Heavy Neutrino Masses,” Phys. Rev. Lett. 39, 165 (1977). doi:10.1103/PhysRevLett.39.165

[2] E. Aprile et al. [XENON Collaboration], “Dark Matter Search Results from a One Ton-Year Exposure of XENON1T,” Phys. Rev. Lett. 121, no. 11, 111302 (2018) doi:10.1103/PhysRevLett.121.111302 [arXiv:1805.12562 [astro-ph.CO]].

[3] K. Griest and M. Kamionkowski, “Unitarity Limits on the Mass and Radius of Dark Matter Particles,” Phys. Rev. Lett. 64, 615 (1990). doi:10.1103/PhysRevLett.64.615

[4] J. Smirnov and J. F. Beacom, “TeV-Scale Thermal WIMPs: Unitarity and its Consequences,” Phys. Rev. D 100, no. 4, 043029 (2019) doi:10.1103/PhysRevD.100.043029 [arXiv:1904.11503 [hep-ph]].

[5] D. J. H. Chung, E. W. Kolb and A. Riotto, Phys. Rev. Lett. 81, 4048 (1998) doi:10.1103/PhysRevLett.81.4048 [hep-ph/9805473].

[6] D. J. H. Chung, E. W. Kolb, A. Riotto and I. I. Tkachev, Phys. Rev. D 62, 043508 (2000) doi:10.1103/PhysRevD.62.043508 [hep-ph/9910437].

[7] M. A. Fedderke, E. W. Kolb and M. Wyman, Phys. Rev. D 91, no. 6, 063505 (2015) doi:10.1103/PhysRevD.91.063505 [arXiv:1409.1584 [astro-ph.CO]].

[8] D. J. H. Chung, E. W. Kolb and A. Riotto, Phys. Rev. D 59, 023501 (1998) doi:10.1103/PhysRevD.59.023501 [hep-ph/9802238].

[9] V. Kuzmin and I. Tkachev, Phys. Rev. D 59, 123006 (1999) doi:10.1103/PhysRevD.59.123006 [hep-ph/9809547].

[10] E. W. Kolb and A. J. Long, Phys. Rev. D 96, no. 10, 103540 (2017) doi:10.1103/PhysRevD.96.103540 [arXiv:1708.04293 [astro-ph.CO]].

[11] H. Kim and E. Kuflik, Phys. Rev. Lett. 123, no. 19, 191801 (2019) doi:10.1103/PhysRevLett.123.191801 [arXiv:1906.00981 [hep-ph]].

[12] T. Hambye, A. Strumia and D. Teresi, “Supercool Dark Matter,” JHEP 1808 (2018) 188 doi:10.1007/JHEP08(2018)188 [arXiv:1805.01473 [hep-ph]].

[13] P. Baratella, A. Pomarol and F. Rompineve, “The Supercooled Universe,” JHEP 1903, 100 (2019) doi:10.1007/JHEP03(2019)100 [arXiv:1812.06996 [hep-ph]].

[14] J. R. Espinosa, T. Konstandin, J. M. No and G. Servant, “Energy Budget of Cosmological First-order Phase Transitions,” JCAP 1006, 028 (2010) doi:10.1088/1475-7516/2010/06/028 [arXiv:1004.4187 [hep-ph]].

[15] D. Bodeker and G. D. Moore, JCAP 0905, 009 (2009) doi:10.1088/1475-7516/2009/05/009 [arXiv:0903.4099 [hep-ph]].

[16] D. Bodeker and G. D. Moore, JCAP 1705, 025 (2017) doi:10.1088/1475-7516/2017/05/025 [arXiv:1703.08215 [hep-ph]].

[17] A. Falkowski and J. M. No, JHEP 1302, 034 (2013) doi:10.1007/JHEP02(2013)034 [arXiv:1211.5615 [hep-ph]].

[18] P. Gondolo and G. Gelmini, Nucl. Phys. B 360, 145 (1991). doi:10.1016/0550-3213(91)90438-4

[19] A. Kosowsky, M. S. Turner and R. Watkins, “Gravitational radiation from colliding vacuum bubbles,” Phys. Rev. D 45, 4514 (1992). doi:10.1103/PhysRevD.45.4514

[20] A. Kosowsky, M. S. Turner and R. Watkins, “Gravitational waves from first order cosmological phase transitions,” Phys. Rev. Lett. 69, 2026 (1992). doi:10.1103/PhysRevLett.69.2026

[21] A. Kosowsky and M. S. Turner, “Gravitational radiation from colliding vacuum bubbles: envelope approximation to many bubble collisions,” Phys. Rev. D 47, 4372 (1993) doi:10.1103/PhysRevD.47.4372 [astro-ph/9211004].

[22] M. Kamionkowski, A. Kosowsky and M. S. Turner, “Gravitational radiation from first order phase transitions,” Phys. Rev. D 49, 2837 (1994) doi:10.1103/PhysRevD.49.2837 [astro-ph/9310044].

[23] D. Cutting, M. Hindmarsh and D. J. Weir, Phys. Rev. D 97, no. 12, 123513 (2018) doi:10.1103/PhysRevD.97.123513 [arXiv:1802.05712 [astro-ph.CO]].

[24] J. Ellis, M. Lewicki, J. M. No and V. Vaskonen, “Gravitational wave energy budget in strongly supercooled phase transitions,” JCAP 1903, 024 (2019) doi:10.1088/1475-7516/2019/03/024 [arXiv:1903.09642 [hep-ph]].

[25] C. Caprini et al., “Detecting gravitational waves from cosmological phase transitions with LISA: an update,” arXiv:1910.13125 [astro-ph.CO].

[26] P. Amaro-Seoane et al. [LISA Collaboration], “Laser Interferometer Space Antenna,” arXiv:1702.00786 [astro-ph.IM].

[27] P. W. Graham, J. M. Hogan, M. A. Kasevich and P. Wolf, Phys. Rev. D 94, no. 10, 104022 (2016) doi:10.1103/PhysRevD.94.104022 [arXiv:1606.01860 [physics.atom-ph]].

[28] P. W. Graham et al. [MAGIS Collaboration], arXiv:1711.02225 [astro-ph.IM].
[29] M. Punturo et al., Class. Quant. Grav. 27, 194002 (2010). doi:10.1088/0264-9381/27/19/194002

[30] S. Hild et al., Class. Quant. Grav. 28, 094013 (2011) doi:10.1088/0264-9381/28/9/094013 [arXiv:1012.0908 [gr-qc]].

[31] S. Kawamura et al., Class. Quant. Grav. 23, S125 (2006). doi:10.1088/0264-9381/23/8/S17

[32] K. Yagi and N. Seto, Phys. Rev. D 83, 044011 (2011) Erratum: [Phys. Rev. D 95, no. 10, 109901 (2017)] doi:10.1103/PhysRevD.95.109901, 10.1103/PhysRevD.83.044011 [arXiv:1101.3940 [astro-ph.CO]].

[33] N. Arkani-Hamed, G. F. Giudice, M. A. Luty and R. Rattazzi, Phys. Rev. D 58, 115005 (1998) doi:10.1103/PhysRevD.58.115005 [hep-ph/9803290].

[34] T. Asaka and M. Yamaguchi, Phys. Lett. B 437, 51 (1998) doi:10.1016/S0370-2693(98)00890-9 [hep-ph/9805449].

[35] T. Asaka and M. Yamaguchi, Phys. Rev. D 59, 125003 (1999) doi:10.1103/PhysRevD.59.125003 [hep-ph/9911451].

[36] N. Abe, T. Moroi and M. Yamaguchi, JHEP 0201, 010 (2002) doi:10.1088/1126-6708/2002/01/010 [hep-ph/0111155].

[37] S. Nakamura, K. i. Okumura and M. Yamaguchi, Phys. Rev. D 77, 115027 (2008) doi:10.1103/PhysRevD.77.115027 [arXiv:0803.3725 [hep-ph]].

[38] K. Choi, E. J. Chun, H. D. Kim, W. I. Park and C. S. Shin, Phys. Rev. D 83, 123503 (2011) doi:10.1103/PhysRevD.83.123503 [arXiv:1102.2900 [hep-ph]].

[39] K. S. Jeong and M. Yamaguchi, JHEP 1107, 124 (2011) doi:10.1007/JHEP07(2011)124 [arXiv:1102.3301 [hep-ph]].

[40] K. Nakayama and N. Yokozaki, JHEP 1211, 158 (2012) doi:10.1007/JHEP11(2012)158 [arXiv:1204.5420 [hep-ph]].

[41] T. Moroi, K. Mukaida, K. Nakayama and M. Takimoto, JHEP 1306, 040 (2013) doi:10.1007/JHEP06(2013)040 [arXiv:1304.6597 [hep-ph]].

[42] R. Jinno, T. Konstandin and M. Takimoto, “Relativistic bubble collisionsa closer look,” JCAP 1909, 035 (2019) doi:10.1088/1475-7516/2019/09/035 [arXiv:1906.02588 [hep-ph]].

[43] M. J. Baker, J. Kopp and A. J. Long, arXiv:1912.02830 [hep-ph].